The Production of Hypernuclei - A Strange Processes

E. V. Hungerford

Department of Physics
University of Houston
4800 Calhoun Rd
Houston, TX 77204, USA

1. - Introduction

The golden anniversary of hypernuclear physics passed several years ago with little celebration. However, this approximately 50 year-old discovery [1] was a seminal step forward in the development of hadronic physics [2]. Recall that in the early 1950's, the hypercharge quantum number, and even parity non-conservation, were unknown. Into this world the concept of SU(3) flavor symmetry was born, providing the symmetry foundation for the present flavor structure of physics.

Early studies of strange nuclear systems used cosmic rays and nuclear emulsion, to insert, and observe, a hyperon within an atomic nucleus, e.g. figure 1. By 1956 a sample of 72 Λ hyperfragments had been collected[3] and studied. This research eventually led to the reported observation of Ξ, and even double Λ hypernuclei, but formation rates of these multistrange systems were limited, and left unverified. Almost all information, even at this date, is limited to single Λ hypernuclei.

Over subsequent years, progress in hypernuclear physics has been driven by the availability of new technologies, developing with brief bursts of energy as new beams and detectors became available. Presently there appears to be a rekindled interest in this field due to the advent of new accelerators, providing intense beams of kaons, pions, anti-protons, and electrons. Therefore, this review will consider some of the general features of cold, strange, hadronic systems. There have been several good reviews over the years [4, 5, 6], however, because of the close coupling between beams, production mecha-
nisms, and detectors, it will be this author’s intention to explore these issues coherently. For example, a study of a particular production mechanism may be optimized by using a particular detector, and a selective application of a specific reaction could result in the emphasis of a particular physics. Of course what can be developed here is only a set of examples. The application of new technologies are left, as always, to the skill and imagination of those working in the field.

2. – Background

A nuclear system may contain one or more bound hyperons. This system is called a hypernucleus, and the investigation of strangeness in the nucleus has proved to not merely be an extension of conventional nuclear physics. The discussion in the present section will mainly develop the physics of single \( \Lambda \) hypernuclei as an example, to establish a foundation, laying out the general features and characteristics of strangeness production, before details of specific production mechanisms are discussed. In any event, only one case of a bound, \( \Sigma \) hypernuclear system has been observed, and evidence for bound, multi-strange systems is meager and controversial. Still bound, multi-\( \Lambda \) nuclear systems must certainly exist.

It is difficult to experimentally determine the \( \Lambda N \) interaction by \( \Lambda \) scattering, so the elementary interactions of a \( \Lambda \) with a nucleon are poorly known. However, hypernuclear physics allows the study of the \( SU(3)_{flavor} \) baryon-baryon interaction at normal nuclear
densities, as an effective $\Lambda N$ potential can be extracted from hypernuclear spectra[7]. This effective potential can then be compared to the form obtained of an elementary interaction, when this form is inserted into the hadronic many-body problem.

The effective potential can then serve as a normalization point, to extrapolate the $SU(3)_{flavor}$ interaction to matter-densities found in neutron stars, where mixtures of nucleons and hyperons could form a stable system[8, 9], figure 2. At higher temperatures, the $\Lambda N$ interaction is also relevant to the cooling of the particle “plasma” formed in relativistic heavy ion collisions, i.e. coalescence.

Of more conventional interest, the strangeness degree of freedom allows the nucleus to rearrange, taking advantage of SU(3) flavor symmetry, to maximize the nuclear binding energy[11]. However, the one pion exchange (OPE) between a $\Lambda$ and a nucleon does not occur due to conservation of isospin. Thus the interaction has shorter range than OPE, and is dominated by higher mass meson exchanges, figure 3, and the two-pion exchange coupling of a lambda to a nucleon through an intermediate sigma ($\Lambda N \rightarrow \Sigma N \rightarrow \Lambda N$). This leads to sizable charge asymmetry and three-body forces[12]. Here in comparison, one needs to not only recall that the $NN$ interaction is dominated by OPE, but that second order diagrams involving an intermediate $\Sigma$ state are more important than similar diagrams involving a $\Delta$, as the $\Sigma$-$\Lambda$ mass difference is much smaller than that of the $\Delta$-$N$, and of course, the $\Sigma$ is much narrower in width.

Finally the $\Lambda$ can be used as a probe of the nuclear medium. If the $\Lambda$ is considered a fundamental particle and remains identifiable as such within the nucleus, a hypernuclear
\( \Lambda \) will sample the nuclear interior where there is little direct information on the single particle structure of nuclei. Of course any system can be expanded in a complete set of wave functions which span this space, so the question really becomes; “how well does the single particle picture describe the interior of a heavy nucleus?” Of relevance to this question, various features of hypernuclei such as \( \Lambda \) decay, and the spectra of heavier hypernuclear systems can be extremely interesting. Thus hypernuclei can better illuminate features which would be more obscured in conventional nuclear systems. Certainly one cannot, nor would one want to, reproduce the wealth of information that has been accumulated on conventional nuclei. However, the hypernucleus offers a selective probe of the hadronic many-body problem, providing insight in areas that cannot be easily addressed by other techniques.

21. In the Beginning. – In the beginning, the production of hypernuclei was limited by cosmic ray interactions in emulsion, for reference see figure 1. As an example, the reaction, \( K_{slp} + A \rightarrow \pi + \gamma A' + X \), produces the strange nuclear system, \( \gamma A' \). This system can then be detected in nuclear emulsion by studying the incident and decay tracks created in the reaction. Although cosmic rays were later replaced by protons from terrestrial accelerators, emulsion or liquid bubble chambers coupled with stopping \( K^- \) beams, remained the detectors and beams of choice. These essentially \( K^- \) spallation interactions with emulsion atoms, limited production to light hypernuclei where sufficient recoil momentum could be observed and measured. While emulsion provided good energy resolution and complete kinematic reconstruction of a hypernuclear decay, scanning was tedious, atomic species limited to those produced from emulsion (or bubble chamber) targets, and statistics were poor. Still the masses of many s- and p-shell hypernuclei were determined by this technique. Table I lists observed species of \( \Lambda \) hypernuclei where the \( \Lambda \) binding energy has been determined[13]. One sees more easily from figure 4 that saturation of \( B_\Lambda \) occurs as \( \Lambda \) increases.

The free \( \Lambda \) lifetime is approximately 0.26 ns, and is mediated by the weak decay of the strange quark. This is much longer than the time scale of hadronic, and most electromagnetic, interactions. Thus \( \Lambda \) weak decay does not effect the \( \Lambda \)-Nuclear many-body system. However, it was quickly recognized that the weak decay of a hypernucleus was unique. Although a free \( \Lambda \) decays by meson emission, a nuclear \( \Lambda \) finds this channel increasingly closed as the atomic number increases. This occurs because as the \( \Lambda \) becomes more bound, the recoil nucleon becomes more Pauli blocked, and in heavy hypernuclei,
mesonic decay is suppressed by some 2 orders of magnitude  [14]. That this suppression seems to saturate at about 1% of the free lifetime, is due to high momentum components and correlations in the single particle wave functions, providing an interesting window on these little measured nuclear features.

For hypernuclei having $A > 6$, hypernuclear decay predominately occurs via the non-mesonic process, $\Lambda N \rightarrow NN$. This is a weak interaction, having a 4-fermion vertex, and is the only observed reaction of this type. The process occurs because the emitted nucleons have momentum higher than the Fermi-surface. As it turns out, mesonic decay decreases and non-mesonic decay increases as the atomic number of the nucleus increases. However the hypernuclear lifetime remains more-or-less constant, decreasing by perhaps a factor of 2 when it saturates in heavier systems [4, 14].

3. – Hypernuclear Structure

Analysis of hypernuclear spectra begins with a potential form for the $\Lambda N$ interaction. This is then folded into the hadronic many-body problem to obtain an effective single particle potential for a $\Lambda$ within the nucleus. Thus the hypernucleus can be described by a set of single particle basis states  [7, 11]. The assumption is that a nuclear $\Lambda$ can be well represented by a superposition of a few single particle states, interacting via an effective $\Lambda N$ potential within the nuclear medium.

However, the direct use of $\Lambda N$ data to obtain the “elementary” $\Lambda N$ interaction potential is severely restricted due to the difficulty in making direct measurements of $\Lambda N$ scattering. There are only $\approx 600$, low momentum (200-300 MeV/c) $\Lambda N$ and $\Sigma N$ scattering events. These data even fail to define the relative sizes of the spin-triplet and
Table I. – Experimental Λ Binding Energies of Hypernuclei

<table>
<thead>
<tr>
<th>Hypernucleus</th>
<th>$B_\Lambda$</th>
<th>$B_\Lambda/\Lambda$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^3_\Lambda$H</td>
<td>0.13</td>
<td>0.943</td>
</tr>
<tr>
<td>$^4_\Lambda$He</td>
<td>2.04</td>
<td>0.510</td>
</tr>
<tr>
<td>$^5_\Lambda$He</td>
<td>2.39</td>
<td>0.598</td>
</tr>
<tr>
<td>$^6_\Lambda$He</td>
<td>3.12</td>
<td>0.624</td>
</tr>
<tr>
<td>$^8_\Lambda$He</td>
<td>4.18</td>
<td>0.697</td>
</tr>
<tr>
<td>$^9_\Lambda$Li</td>
<td>7.16</td>
<td>0.563</td>
</tr>
<tr>
<td>$^{11}_\Lambda$Li</td>
<td>4.50</td>
<td>0.750</td>
</tr>
<tr>
<td>$^{11}_\Lambda$Be</td>
<td>5.58</td>
<td>0.797</td>
</tr>
<tr>
<td>$^{12}_\Lambda$Be</td>
<td>6.80</td>
<td>0.850</td>
</tr>
<tr>
<td>$^{12}_\Lambda$Be</td>
<td>6.71</td>
<td>0.746</td>
</tr>
<tr>
<td>$^{13}_\Lambda$Be</td>
<td>9.11</td>
<td>0.911</td>
</tr>
<tr>
<td>$^{14}_\Lambda$Be</td>
<td>8.29</td>
<td>0.921</td>
</tr>
<tr>
<td>$^{15}_\Lambda$B</td>
<td>8.89</td>
<td>0.889</td>
</tr>
<tr>
<td>$^{16}_\Lambda$B</td>
<td>10.24</td>
<td>0.931</td>
</tr>
<tr>
<td>$^{17}_\Lambda$B</td>
<td>11.37</td>
<td>0.948</td>
</tr>
<tr>
<td>$^{12}_\Lambda$C</td>
<td>10.80</td>
<td>0.900</td>
</tr>
<tr>
<td>$^{13}_\Lambda$C</td>
<td>11.69</td>
<td>0.899</td>
</tr>
<tr>
<td>$^{14}_\Lambda$C</td>
<td>12.17</td>
<td>0.869</td>
</tr>
<tr>
<td>$^{15}_\Lambda$N</td>
<td>12.17</td>
<td>0.809</td>
</tr>
<tr>
<td>$^{16}_\Lambda$N</td>
<td>13.59</td>
<td>0.906</td>
</tr>
<tr>
<td>$^{17}_\Lambda$O</td>
<td>12.50</td>
<td>0.781</td>
</tr>
<tr>
<td>$^{28}_\Lambda$Si</td>
<td>16.6</td>
<td>0.592</td>
</tr>
<tr>
<td>$^{32}_\Lambda$Si</td>
<td>17.5</td>
<td>0.547</td>
</tr>
<tr>
<td>$^{40}_\Lambda$Ca</td>
<td>20.0</td>
<td>0.500</td>
</tr>
<tr>
<td>$^{41}_\Lambda$V</td>
<td>19.5</td>
<td>0.382</td>
</tr>
<tr>
<td>$^{56}_\Lambda$Fe</td>
<td>21.0</td>
<td>0.375</td>
</tr>
<tr>
<td>$^{90}_\Lambda$Y</td>
<td>23.2</td>
<td>0.261</td>
</tr>
<tr>
<td>$^{139}_\Lambda$La</td>
<td>23.8</td>
<td>0.171</td>
</tr>
<tr>
<td>$^{208}_\Lambda$Pb</td>
<td>26.5</td>
<td>0.127</td>
</tr>
</tbody>
</table>

spin-singlet, s-wave scattering lengths and effective ranges, and a phase shift analysis of these data is not practical [5].

Progress has been made by combining the ΛN data with the much more copious NN data, using $SU(3)_f$ flavor symmetry [15]. This analysis produces general baryon-baryon scattering amplitudes, and several potential forms have been used to represent the data. The Nijmegen D, and F, or the newer, soft-core forms, NSC97 and 99, are the most predominately used potentials [16, 17]. Obviously $SU(3)_f$ is badly broken due to the difference between the s and (u,d) quark masses, but the potentials attempt to account for this mass difference by adding mass breaking terms. Table II provides several simple
comparisons for various ΛN potentials.

**Table II. – Comparison of various model ΛN amplitudes**

<table>
<thead>
<tr>
<th>Model</th>
<th>Ref.</th>
<th>$a^t$</th>
<th>$r_0^t$</th>
<th>$a^l$</th>
<th>$r_0^l$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Nijmegen D</td>
<td>[22]</td>
<td>-1.90</td>
<td>3.72</td>
<td>-1.96</td>
<td>3.24</td>
</tr>
<tr>
<td>Nijmegen F</td>
<td>[22]</td>
<td>-2.29</td>
<td>3.17</td>
<td>-1.88</td>
<td>3.36</td>
</tr>
<tr>
<td>Nijmegen SC</td>
<td>[22]</td>
<td>-2.78</td>
<td>2.88</td>
<td>-1.41</td>
<td>3.11</td>
</tr>
<tr>
<td>Julich A</td>
<td>[22]</td>
<td>-1.56</td>
<td>1.43</td>
<td>-1.59</td>
<td>3.16</td>
</tr>
</tbody>
</table>

Within a nucleus the general hyperon-nucleon potential can be expressed by the form;

\[
V(r) = V_0(r) + V_s(S_N \cdot S_Y) + V_{12} + V_{ls}(L \times S^+) + V_{al}(L \times S^-). \tag{1a}
\]

In this expression $S_{12} = 3(\sigma_1 \cdot \hat{r})(\sigma_2 \cdot \hat{r}) - \sigma_1 \cdot \sigma_2$ is the usual spin-tensor operator, and $S^\pm = 1/2(S_N \pm S_Y)$ are symmetric and anti-symmetric combinations of nucleon and hyperon spin operators. The new feature in this effective potential is the anti-symmetric spin-orbit operator which vanishes for the NN system. The net spin-orbit is obtained by combining the $V_{ls}$ and $V_{al}$ terms. In meson exchange potentials, both of these terms are small [4, 18]. In potentials derived from interacting quarks, the terms can be large, but have opposite signs resulting in a small, nearly zero, spin orbit potential [19]. Both meson exchange and quark potentials reproduce the small size of the measured spin orbit term.

The single particle potential form as illustrated above, can be used to construct the hypernuclear wave functions, which are then used to calculate wave functions for the transition matrices of the production reactions.

3'1. Comparison of the ΛN and ΣN Potentials. – The Σ is an isospin 1 particle, so that its interaction with a nucleon can occur in either an isospin 3/2 or 1/2 state. This is opposed to the isospin 1/2 state obtained by coupling the isospin zero Λ to a nucleon. The ΣN potential can take the same form as equation 1a, with the exception that the $V_i$ are isospin dependent [20]. The Nijmegen potentials indicate, and the data are consistent with, significant spin dependence, predicting strong attraction in the $^1S_0$, $T = 3/2$ and $^3S_1$, $T = 1/2$ ΣN channels and weak repulsion in the $^3S_1$, $I = 3/2$, and $^1S_0$, $T = 1/2$ channels. Of most importance, however, is the absorptive channel, ΣN → ΛN, which introduces imaginary terms in the ΣN potential. The absorptive behavior of the ΣN potential is large and dominates the behavior of Σ hypernuclei.

The strong coupling, ΛN → ΣN, also influences the ΛN potential, as the Λ-Σ nuclear system is really a coupled channel problem, where virtual Σ interactions are dependent on the spin and isospin of a particular hypernuclear level. This has been shown, for example, to significantly effect the level structure of s-shell hypernuclei [21] where multibody calculations of the structure can be undertaken figure 5. Note that there are four
s-shell hypernuclei, and the 5-body system is bound. Thus the $\Lambda$ as a distinguishable particle, resides in the s-shell and adds binding to the system. Gammas from the excited states in the 4-body system have been observed.

4. The Production Formalism

In order to discuss hypernuclear production, figure 6 provides a cartoon of quark flows for several reaction processes. In each case a strange quark is deposited in the nucleus either by exchange or by associative production. This can be achieved by a number of reactions, and the figure does not provide quantitative detail, but it does provide a pictorial display of the process. In this section the general features of hypernuclear production, and the resulting spectrum, will be introduced, while specific reactions will be discussed in later sections.

It has been found that hypernuclear production can be described by the distorted wave impulse approximation [22], DWIA. In its simplest form, DWIA is applied without inclusion of spin, neglecting multi-step processes (such as intermediate $\Sigma$ production), and off-shell effects. Spin, in particular, can be added with some additional complexity. However, in order to apply DWIA, one must have knowledge of the elementary amplitudes for hyperon production, which in general are not well known. In addition, the distorted waves are calculated from an optical potential involving the initial and final states, so this information is also required. The DWIA form for a reaction such as $K + A \rightarrow \pi + \gamma A$ is;
\[
\begin{align*}
\gamma + p &\rightarrow \Lambda + K^+ \\
\pi^+ + n &\rightarrow \Lambda + K^+
\end{align*}
\]

Fig. 6. – Quark flow diagrams for various production mechanisms

\[T_{AY} \approx < \chi_\pi^- | T | \chi_K^+ \psi_A > \approx < \pi \Lambda | t | K N > N_{eff}^{1/2}\]

In this equation the \(\chi^\pm\) are the distorted incident and final wave functions for the kaon and pion in the nuclear optical potentials, and \(t\) is the elementary transition matrix. The final form has been factored so that the transition amplitude, \(t\), multiplies a density representing the effective number of nucleons participating in the interaction. Following [22] this density is written:

\[N_{eff}^{1/2} = \int d^3x \ < \chi_\pi^- | x > < x | \chi_K^+ > < \psi_Y | \Lambda^+ N | 0 >\]

The \(< \psi_Y | \Lambda^+ N | 0 >\) matrix element in the above equation accounts for the \(N \rightarrow \Lambda\) transition, while the distorted waves are normalized so that \(N_{eff} = 1\) for the elementary process. The distorted waves can be expanded in spherical harmonics about the incident beam direction, giving a spherical tensor dependent on the angular momentum transfer, \(\Delta l\). For spinless transitions, only natural parity excitations are possible, and for small momentum transfers, \(\Delta l = 0\). One finds that the results are not so sensitive to the details of the elementary reaction, at least at the present level of experimental accuracy.

Distortions of the initial and final wave functions generally do not change the shape of the differential cross sections, but can decrease their size by an order of magnitude.
Absorption influences which states which are produced, as for example, strong absorption restricts the reaction to occur at the nuclear surface. This may be seen in the simplest eikonal approximation to the effective nucleon number [4], which is given by;

\[
N_{\text{eff}} = \int d^3r \, \rho_n(r) \exp\left[-\sigma_k \int_r^\infty \rho \, dz' - \sigma_{\text{pt}} \int_z^\infty \rho \, dz'\right].
\]

(4a)

In the above expression, \( \rho_n \) is the neutron density and \( \rho \) the total nuclear density. The eikonal approximation is strictly valid for zero degree scattering, but it clearly shows the effect of absorption.

The DWIA approximates the production process, assuming that an incident projectile interacts with a nucleon in a specific, single-particle nuclear state, producing a hyperon which is then deposited in a specific, single-particle hyperon state. More complexity can be incorporated by superimposing additional single particle components into the wave functions. However in this formalism, both the production and the spectra of hypernuclei can be considered in terms of [particle, hole] states. Here a “hole” represents the removal of a nucleon from a nuclear wavefunction and representing the “hole” as the wavefunction of the residual nuclear core. The “hole” has the quantum numbers of the nucleon which was removed. A hyperon single particle state is coupled to this “hole” to produce the hypernuclear state, figure 7. The procedure is simplest for nuclei near closed shells.

When viewed in this way, one immediately understands why kinematics strongly influences hypernuclear production. figure 8 gives a comparison of the \( \Lambda \) recoil momentum as a function of the incident projectile momentum. If the momentum transfer is above the Fermi-momentum, the \( \Lambda \) can escape the nucleus, leading to continuum, rather than bound, states. The low linear momentum transfer reaction, \((K^-, \pi^-)\), is a “substitution” reaction as it replaces a nucleon with a hyperon having the same spin and angular momentum quantum numbers, since the angular momentum transfer, \( \Delta l \approx 0 \). For higher momentum transfers, angular momentum matching leads to higher production cross sections, and the \((\pi^+, K^+)\) reaction preferentially excites “spin-stretched” states [23].

Because the nucleons on which the reaction takes place, move within the nuclear potential well, and the elementary cross section is factored in the “\( t \rho \)” approximation, the elementary \( t \) matrix amplitude should be averaged over the nuclear Fermi momentum. This lowers the cross section by 10-20\%. There are several ways to do the averaging, but the result is not sensitive to the method employed. figure 10 shows the result of Fermi averaging the elementary amplitudes or the cross sections for the \((K^-, \pi^-)\) reaction [24].

Finally DWIA assumes that the reaction occurs via interactions which are contained within the elementary, on shell \( t \) matrix. Of course this amplitude should really be extrapolated off-shell, and it also does not contain the full range of interaction possibilities when the interaction takes place in a nuclear environment. The error associated with using on-shell values is estimated [22] to be on the order of 30\%. Exclusion of reaction
Fig. 7. – An example of the extreme (particle, hole) model for the hypernucleus. In this example, an s-shell $\Lambda$ is added to the $^{11}$C nuclear core after removal of a neutron from the $^{12}$C nucleus.

Fig. 8. – The recoil momentum of the hyperon in various elementary production reactions at $0^\circ$ as a function of the incident momentum.
processes involving more than one nucleon, such as the production of intermediate $\Sigma$ states, are estimated to be small [22].

4'1. Production of Continuum States. – In many situations the $\Lambda$ is produced in unbound, continuum states. This especially occurs in high momentum transfer reactions, but even in the $(K^-, \pi^-)$ substitutional reaction $\Lambda$ can be unbound, as the $\Lambda$-Nucleus well depth is approximately 1/2 of that of the nucleon-Nucleus well depth. In hypernuclear production, this leads to the creation of a continuum background of excitations above the $\Lambda$-Nucleus threshold.

The continuum is sometimes discussed in terms of a “quasi-free”, QF, reaction. In
this model, the QF continuum spectrum is obtained by calculating the statistical density of states for the reaction on a single-particle nuclear state which produces an unbound $\Lambda$ recoiling under the influence of a $\Lambda$-Nucleus potential. Calculations of the spectrum are usually undertaken in a Fermi-gas model, so that the shape of the spectrum is determined by kinematics and the $\Lambda$-Nucleus well depth [25].

In any 3 body reaction process, energy and momentum conservation can be written;

\begin{align}
\omega + M_A &= E_2 + E_3 \\
\vec{q} &= \vec{p}_2 + \vec{p}_3;
\end{align}

where $\omega = E_0 - E_1$ is the energy transfer, and $\vec{q} = \vec{p}_0 - \vec{p}_1$ is the 3-momentum transfer. One then models the nucleus, $A$, as a core plus a bound nucleon, $M_A = M_{A-1} + M_N + B_N$, where $B_N$ is the binding energy. In a Fermi-gas model this is related to the well depth of the nucleon potential. For example, in a reaction such as, $\pi + A \to K + \Lambda + (A - 1)$, particle 0 would be the $\pi$, particle 1 the kaon, particle 2 the $\Lambda$ and particle 3 the recoiling core, $(A - 1)$. The core remains a spectator to the reaction and is assumed to recoil with the Fermi-momentum that it had before the reaction, $-\vec{p}_0$. The recoiling core, and generally the $\Lambda$ as well, can be treated non-relativistically. The statistical density of states for the kaon momentum is then;

\begin{align}
N \sim \int d^3p_2 \, d^3p_3 \, \delta(\vec{q} - \vec{p}_2 - \vec{p}_3) \delta(\omega - E_2 - E_3).
\end{align}

The QF spectrum shape is determined by solving the kinematic equations and performing the integrals. Making a reasonable approximation [25] in order to obtain an analytic solution, a parabolic spectrum shape is found having a maximum at $\varpi$ given by;

\begin{align}
\varpi = M_\Lambda - M_N + (U_N - U_\Lambda) - (M_\Lambda - M_N)k^2/(4M_NM_N) + q^2/(2M_\Lambda).
\end{align}

Here, $M_\Lambda$ and $M_N$ are the $\Lambda$ and $N$ masses, and $U_\Lambda$ and $U_N$ the well depths. The Fermi momentum is $k_f$ and $\vec{q}$ and $\omega$ are the momentum, $\vec{q} = p_k - \vec{p}$, and energy, $\omega = E_K - E_\pi$ transfers. Applying this analysis to the continuum data of several medium $\Lambda$ hypernuclei, a $\Lambda$-Nucleus well depth of $\approx 30$ Mev is extracted, figure 11.

On the other hand, contributions to the continuum spectrum should also include resonant behavior, i.e. nuclear structure information. Inclusion of nuclear structure can be treated by several methods [26, 27], the most common being the continuum shell model [28], where the QF and resonant behavior are simultaneously calculated. The general features of the continuum production are best observed by comparing the spectra from various reactions [4], figure 12. This will be discussed in more detail later, but in the
case of \((K^- , \pi^- )\), states near the surface are excited and the QF component is small. This is due to the selective nature of the reaction and the low momentum transfer. In high momentum transfer reactions such as \((\pi^+ , K^+)\), the quasi-free contribution is larger. However, selectivity in the reaction process can also constrain continuum production.

When the QF process is applied to the \(\Sigma\) hypernuclear continuum, one extracts a real well depth of about 32 MeV, which is similar to that obtained from \(\Sigma^-\) atom data. However this is inconsistent with the well depth extracted using continuum shell model calculations, 5-10 MeV. One finds that the QF approximation does not vanish at particle thresholds, but continues into the unphysical region. The continuum shell model modifies the shape of the spectrum, requiring that the spectrum vanish at thresholds, and includes resonance behavior.

4.2. The Nuclear Auger Effect. - From previous arguments, one can model a hypernucleus as a set of single particle nucleon holes and \(\Lambda\) states. A reaction then can place a \(\Lambda\) particle in any of the bound or unbound levels of the nucleus, from which it either escapes the nuclear potential well or cascades downward in energy [29], eventually reaching the ground state before it weak decays, unless it becomes trapped in an isomeric level, figure 13. The energy released in these transitions can be removed either by gamma rays or Auger neutron emission. In a hypernucleus, the neutron emission threshold can be lower than the \(\Lambda\) threshold, and in any event, nucleon emission can occur even from unbound \(\Lambda\) states. Thus the final hypernuclear species may differ from the one initially produced. Indeed in the hypernuclear system may fission, producing a hypernucleus much lower in mass.

This effect is illustrated in the data [30] from the reaction \(^7\text{Li}(K^- ; \pi^- , \gamma)X_Y\). A pion spectrum of this reaction in terms of the \(B_\Lambda\) value for the hypernucleus \(^3\Lambda\text{Li}\) is shown in figure 14. Energy resolution is poor due to target thickness which was required to obtain a reasonable rate for the coincidence experiment. However there are several regions into which \(B_\Lambda\) can be divided; a) a region encompassing bound states; b) a region of QF production; and c) a region high in the continuum containing an \((s\)-shell neutron hole, \(p\)-shell \(\Lambda\)). Placing cuts on the events produced within these various regions of \(B_\Lambda\), gamma
de-excitations are observed in coincidence. These correspond to; a) $^7_\Lambda$Li hypernuclear transitions, b) $^6$Li nuclear transitions, and c) $^4_\Lambda$H hypernuclear transitions, indicating the system fissions into clusters.

This illustrates that hypernuclei can be studied either by either; 1) production mechanisms where the reaction is constrained to a few measured particles which completely determine the products; or by 2) decay mechanisms where the production process may be ill determined, but measurement of the decay products for a specific hypernucleus are sufficient for its identification. Note however that the hypernucleus in this example was identified because the energy of the gamma was known. Unless some additional
information is available, just measuring gamma energies is not sufficient to identify a hypernucleus or hypernuclear event.

5. **Beams and Production Reactions**

A hyperon, and in particular a hypernucleus, is created by replacing an up or down quark with a strange one. Such reactions involve strangeness exchange or associated production, and quark flow diagrams of these processes are shown in figure 6. Thus in general, one must begin with a flux of particles that contain strange quarks, or a flux with sufficient energy to make an $s\pi$ pair. Of course the intensity of the flux must be sufficient to overcome low production probabilities, and the whole process must be amendable to some type of detection apparatus. If the hypernucleus is to be studied through its production products, the momentum of both the beam and the reaction particles are normally measured using magnetic spectrometers.

Discussions in this section consider only the use of secondary beams used to produce strange nuclear systems. Direct use of a primary beams for this purpose will be discussed in later sections.

5'1. **Kaon Beams.** - The valence quarks of the $K^-$ are $\bar{u}s$, and the kaon lifetime is about 12 ns, with $ct = 3.7m$. Thus kaons, with sufficient relativistic boost, have a reasonable probability of survival during the traversal of a beamline of 15-20m. They
Fig. 14. – The $^7$Li(K$^-$; $\pi^-$, $\gamma$)$_X$Y spectrum showing the $B_\Lambda$ Spectrum and the Gamma transitions from the various regions of this spectrum.
must be produced as a secondary beam from a target, usually Pt or W, by the interaction of a primary beam of protons. Secondary particle production at zero degrees [31] as a function of beam energy is shown in figure 8. Usually the secondary target is several interaction lengths thick, and acts as a source for the magnetic transport from the production target to the experimental hall.

Kaon production in the target is driven by associated production. The elementary reaction has the form, \( pN \rightarrow KYN' \), where \( Y \) is either a \( \Lambda \) or a \( \Sigma \), \( N \) and \( N' \) are nucleons or nucleon resonances, and the particles must be chosen to balance charge. A little forethought using charge and strangeness conservation shows that positive, as opposed to negative, production should dominate. Indeed near threshold, \( K^+ \) yields are a factor of 10 above those for \( K^- \). This reduces to \( \approx 3 \) for higher momentum kaons if the primary protons are above 10 GeV/c. Finally, the number of kaons decreases rapidly as the momenta of the produced Kaons decreases. This is modeled in \( pN \) collisions by symmetric backward/forward particle production in the barycentric frame, which is then boosted to the lab frame [32].

Secondary particle production and absorption in the target is usually maximized at zero degrees. The secondary flux emerging from a target is approximately given by the equation:

\[
N_p = N_0 \sigma \frac{\lambda_0 - \lambda_p}{\lambda_0 - \lambda_p} \left[ \exp(-t/\lambda_p) - \exp(-t/\lambda_0) \right]
\]

Here, \( N_p \) is the number of particles emerging from the target, \( N_0 \) is the number of interacting protons, \( \sigma \) is the production cross section per unit length of material, and \( \lambda \) are the absorption lengths in the target. Taking the secondary flux at an angle other than zero, can reduce absorption of these particles due to the transit distance in the target, and under some circumstances this can increase the flux.

Of course kaons are not the only particles produced from the target, and even though the beam is momentum selected, kaons represent only a small fraction of the total particle flux entering the secondary beam line. Early use of kaons beams were limited, not only by intensities, but by the large background of other particles in the beam. For these reasons the experiments used stopping \( K^- \) beams.

In the 1970s, more intense, separated beams of kaons became available [33]. A schematic of a separated beamline is shown in figure 15. The forward dipole/quadrupole magnets collect the beam which is then focused through a velocity separator composed of crossed \( \vec{E} \) and \( \vec{B} \) fields. The figure shows two stages of separation with a sextupole between the separators to correct for higher optical aberrations in the magnetic transport system. Electric fields as high as \( \pm 225 \text{ kv} \) are applied to plates which are separated by 10 cm. The system is enclosed in a pressurized tank containing \( SF_6 \) to limit sparking. The magnetic field is tuned to allow particles of a selected mass to travel without deflection through the separator. A horizontal slit, a few mm high (note the figure appears to separate horizontally but the separation is perpendicular to the dispersion direction), is
positioned to select an enriched beam of a particular mass. This “mass” slit then acts as the object for the magnetic transport into the experimental cave. The scintillation hodoscope, S1, is used in the beam tracking analysis to tag the incident beam momentum on the target.

The performance of a separator depends on the beam momentum, separation length, and optical quality of the beamline. In the case of the beamline described above, one finds a $\pi^-/K^-$ ratio of $(5-10)/1$ for momenta near 800 MeV/c. Similar beamlines operated at higher momenta produce a $\pi/K$ ratio of 1/1. The parameters of a newly designed beamline at the Alternating Gradient Synchrotron, AGS, of the Brookhaven National Laboratory, BNL, are given in Table III. This beamline [34] is designed specifically for stopping $K^+$ in a rare kaon decay experiment, but could satisfy requirements for a new low momentum kaon line for hypernuclear research. Even though separated, backgrounds are still significant, and kaons must be individually tagged. The kaon flux can be as high as $10^6 - 10^7$ s$^{-1}$ during a beam spill, which can last several seconds with perhaps an average of 1000 spills/hr. Such beams have a momentum spread of $\delta p/p \approx 5\%$. Beam transport from the mass slit to the experimental target can be designed to provide a dispersed or non-dispersed focus to match the experimental requirements.

5’2. Momentum Considerations in Applying Kaons to Hypernuclear Production. – Observation of the zero degree $n(K^-, \pi^-)\Lambda$ differential cross section indicates a large enhancement, figure 10, between the momenta of 700 and 900 MeV/c. This enhancement is also near the value producing zero recoil momentum, figure 8. In fact by allowing a small momentum transfer, the central momenta of the forward and backward spectrometers can be different, preventing beam particles from entering the rear spectrometer. There is another enhancement in the elementary cross section at about 1700 MeV/c, but this momentum range has not been explored for $S=-1$ production as it does not seem to offer significant advantages.

In the case of $\Sigma$ production [20], the $N(k^-, \pi)\Sigma$ differential cross section in the forward direction shows two enhancements, figure 16, one at about 400 MeV/c and a smaller one
TABLE III. – *The parameters of the new LEB III beamline at the AGS accellerator.* [34]

<table>
<thead>
<tr>
<th>Particle</th>
<th>Momentum (MeV/c)</th>
<th>Particles/s (a)</th>
<th>Purity</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K^+$</td>
<td>800</td>
<td>$4.8 \times 10^6$</td>
<td>71%</td>
</tr>
<tr>
<td>$K^+$</td>
<td>Stopped</td>
<td>$1.0 \times 10^6$</td>
<td>–</td>
</tr>
</tbody>
</table>

(a) Rate for $10 \times 10^{12}$ particles/s, 25 GeV/c primary beam momentum, 9 cm Pt production target

of different isospin at about 750 MeV/c. The 400 MeV/c momentum is generally too low to be useful as the kaon flux drops rapidly below 600 MeV/c. However zero momentum transfer occurs at an incident momentum of about 300 MeV/c, and QF production is significantly enhanced if an incident momentum value > 600 MeV/c is used.

The production of $S=-2$ hyperons using the $(K^-, K^+)$ reaction requires beam momenta greater than $\approx 1.1$ MeV/c, with a production maximum occurring near 1.5 MeV/c.

![Free space $^7N \to \Sigma$ differential cross sections in lab (0°)](image)

Fig. 16. – The elementary $p(K^-, \pi^-)\Sigma$ production cross section at zero degrees. [4]
At this momenta the kaon flux can be higher, and \(\pi/K\) separation better, but of course double strangeness cross sections are lower, backgrounds higher, and the reaction requires a two-step process.

5.3. Pion Beams. - Secondary pion production is some order of magnitude above kaon production. Again because of charge conservation when using proton beams, \(\pi^+\) is produced more copiously than negative pions by perhaps a factor of 2.5. The pions of interest here, \(>1\text{GeV}/c\), have longer lifetime than kaons, and a higher \(\beta\gamma\tau\) lifetime factor, so there can be a significantly higher flux on target. This can compensate for lower cross sections, but of course the states which are most predominantly produced will also be different.

The \((\pi^+, K^+)\) elementary cross section is shown in figure 17. It peaks at an incident momentum of slightly higher than 1 GeV/c and falls rapidly. In contrast to low momentum kaon induced reactions, the \(\Lambda\) recoil has substantial polarization at forward angles. Unfortunately the maximum in the polarization occurs at the minimum in the production cross section. With the exception that polarization creates specific spin states in the hypernucleus, polarization in the \((\pi^+, K^+)\) reaction has not been experimentally used in spectroscopic studies as production has generally been at very forward angles.

Pion beams will have the same dispersion as kaon beams, \(\delta p/p\) of approximately 5%. At a beam momentum of \(\approx 1\text{GeV}/c\) this means that the uncertainty of the incident beam can dominate the resolution of a spectrometer system. There have been proposals to rotate the phase space of a pulsed pion beam to produce a more intense beam with high resolution [35]. While it is certainly possible to do this, no such plans are presently active.

Thus production of hypernuclei using the elementary process \(n(\pi^+, K^+)Y\) are limited in resolution, as are all secondary beam experiments, by target thickness (statistics) and beam dispersion.

5.4. Anti-proton Beams. - Anti-protons are produced at approximately 1% of the kaon yield. While the anti-protons have long lifetimes in reasonable vacuum, they need to be collected and focused onto a target. Stopping anti-proton beams efficiently use most of the beam in reaction processes, while in-flight reactions are much less efficient. If anti-protons are used for reactions in flight, the momentum width of the beam may limit the experimental resolution, unless they are collected and cooled. However, it is likely that this would not be so important, as the detection of hypernuclear events could rely on the observation of decay products rather than on the reaction products.

6. – Specific Production Reactions for \(S = -1\) Systems

In this section we discuss the production of hypernuclei by several specific reactions. As previously discussed, kinematics play a vital roll in determining the characteristic behavior of a particular production mechanism.
6.1. Hypernuclei Produced by the In-flight, \((K^-, \pi^-)\) Reaction. – The in-flight, \((K^-, \pi^-)\) reaction was not the first used for hypernuclear production, but it was the reaction that introduced the modern era of accelerator based investigations using magnetic spectrometers and electronic counters. This reaction was chosen to be first discussed here, because it provides the most logical introduction to the other reactions.

In the early 1970's it was recognized that the incident momentum in the \(n(K^-, \pi^-)\Lambda\) reaction could be chosen so that the momentum transferred to the \(\Lambda\) could be zero, figure 8. At the same time, separated beams of kaons at approximately 750 MeV/c could provide a maximum in the elementary cross section. Thus a series of experiments using the in-flight \((k^-, \pi^-)\) reaction was initiated at CERN [36] and then at BNL [37]. The spectra of light hypernuclei show peaks for substitutional states (a neutron replaced by a \(\Lambda\) with the same quantum numbers) for states near the nuclear surface. This occurs because the strong nuclear absorption limits \(K^-\) penetration of the nucleus, and zero
angular momentum transfer. Thus as the nuclear radius increases, the excitation of core states decreases. The reaction yield to specific states also decreases, because for a given target thickness in \( \text{gm/cm}^2 \) the number of target nuclei decreases as \( A \) while the cross section increases only as \( A^{2/3} \) resulting in an over all decrease in strength of \( A^{1/3} \).

A measure of the splitting of the p-shell states observed in the \( {}^{16}\text{O}(K^-, \pi^-)\alpha_0\) spectrum shows that the energy differences between the states obtained \[36\] when substituting a \( \Lambda \) for a \( 1/2 \) or \( 3/2 \) neutron is the same as the energy splitting of the hole states in \( {}^{15}\text{O} \). This indicates that the effective \( \Lambda N \) spin-orbit splitting is small. A small effective \( \Lambda N \) spin-orbit potential was also confirmed in the analysis of the angular distribution of the \( {}^{13}\text{C} \) spectrum where the excitation of a p-shell, \( \langle j = 1/2 \text{ neutron hole}, J = 1/2 \text{ or } 3/2 \text{ \( \Lambda \) hypernuclear state depends on the reaction angle \[37\]. Therefore by measuring the shift in this p-shell peak as a function of angle, one can extract the spin-orbit strength.

It is also interesting to note that the high lying \( \Lambda \) states occurring near \( B_\Lambda = 0 \) have narrow width. These states transition to lower energy hypernuclear states by Auger or gamma transitions. Comparable nuclear states would be broad, because they would have deep hole structure. The narrow width of \( \Lambda \)-Nuclear states is due to the weak \( \Lambda \) spin-orbit interaction and the zero isospin \[29\] of the \( \Lambda \).

After the initial success of applying the \((K^-, \pi^-)\) reaction to \( \Lambda \) hypernuclei, an attempt was made to look for bound \( \Sigma \) states \[38\], although it was expected that such structures would have a large width due to the strong conversion, \( \Sigma N \to \Lambda N \). This research remained in a confused state for a number of years, limited by the low statistics of the experiments and speculations by theorists.

Experimentally, a number of light \( \Sigma \)-nuclear systems were investigated, particularly for s- and p-shell, \( \Sigma \) nuclear systems. Attempts were made to use lower incident \( \text{K} \) momentum to reduce the the QF component, and to enhance substitional state production. All systems that were investigated indicated strength below the \( \Lambda \) emission threshold, but the interpretation of the observed structure was limited by statistical fluctuations.

Two high statistics experiments were eventually completed, one \[39\] on a \( \text{\^4He} \) target and one \[40\] on a series of p-shell nuclei. The result provided a consistent picture for \( \Sigma \) nuclear interactions in light nuclear systems. A significant dependence on isospin was found by observing production differences in the spectra from \((K^-, \pi^-)\) and \((K^-, \pi^+)\) reactions. This is shown in figure 18 where one sees a progressive shift of the enhancement below threshold to higher energies and a broadening of its width. In the specific case of \( \text{\^4He} \), a broad bound state having a binding energy of \( \approx 4.4 \text{ MeV} \) with a width of \( \approx 7.0 \text{ MeV} \) was observed. One notes that this state must have isospin \( 1/2 \), as it is seen only in the \((K^-, \pi^-)\) reaction.

The presence of isospin dependence suggests a “Lane” term in the potential which would have a \( 1/A \) dependence, reducing the possibility of \( \Sigma \) states of any width for \( A > 4 \). In particular for the \( A = 4 \) system, theoretical analysis has shown that the effective \( \Sigma \) nucleus potential system has a small attractive pocket near the nuclear surface, with a strong repulsive core and decreasing exponentially as the nuclear radius increases. A bound \( \Sigma \) could reside in this well, and as the nuclear surface has lower density, the conversion width of the \( \Sigma \) decreases, allowing the broad state to form. Given
this understanding, a developed spectroscopy of $\Sigma$ hypernuclei is thus unlikely.

6.2. A Hypernuclei Produced by the $(K^{-}_{\text{stopped}}, \pi^{-})$ Reaction. - Before separated beamlines, the $K^{-}_{\text{stopped}}$ reaction was extensively used to produce hypernuclei. In the first counter experiment of this type [41] a kaon beam was brought to rest in a carbon target, and following the absorption of the kaon, a $^\Lambda_{12}$C hypernucleus was formed and identified by the emission of a $\pi^{-}$. Two broad peaks were observed in the pion spectrum, one with a $B_{\Lambda} = 11 \pm 1$ MeV and the other with $B_{\Lambda} = 0 \pm 1$ MeV. The widths were dominated by the experimental resolution, $6 \pm 1$ MeV/c, and the production rates of the two peaks, subsequently identified as p- and s-shell hypernuclear states in $^\Lambda_{12}$C, were $(2 \pm 1) \times 10^{-4}$, and $(3 \pm 1) \times 10^{-4}$ per stopped kaon, respectively.

In another $K^{-}_{\text{stopped}}$ experiment, the $^{12}\text{C}(K^{-}_{\text{stopped}}, \pi^{-});^{12}\text{C}$ reaction was observed at KEK [42], where it was found that the probabilities per stopped kaon for the formation of s-shell and p-shell states were $(0.98 \pm 0.12) \times 10^{-3}$ and $(2.3 \pm 0.3) \times 10^{-3}$, respectively. These formation probabilities were a factor of 3 larger than the values calculated by Gal [43] and Bando [44], and a factor of 8 larger than the Matsuyama-Yazaki [45] predictions. However, the relative strengths between the two peaks was found to be in better agreement with theory.

The stopped reaction proceeds when a kaon is absorbed from an atomic orbit into the nucleus. X-ray measurements of kaon absorption on $^{12}\text{C}$ [46], for example, indicate that 20% of all the kaons are captured from 3d orbits, while the remaining 80% are believed to be captured from low angular momentum, $l_{\pi}=0$ or 1, and large $n_{\pi}$ states. During capture, a $\Lambda$ is produced by the reaction, $K^{-} + N \to \Lambda + \pi$. Kaon absorption at rest provides momentum transfer approximately equal to the Fermi momentum of a bound
Λ, and for a C target, angular momentum transfers $J \leq 4$ are possible.

More recently a theoretical study of the in-medium modification of the $\vec{K}N$ interaction points out that the $(K^-_{\text{stopped}}, \pi)$ reaction can be used to better define the $K^-$ optical potential at threshold. New estimates of the branching ratios were produced using a $\bar{K}N$ t matrix constructed within a coupled-channel chiral model from the atomic data [47].

As previously discussed, the stopped reaction has higher momentum transfer than the in-flight reaction, and is much less selective, figure 12. In comparison, the QF process is larger so that it becomes difficult to resolve states near $B_{\Lambda} = 0$ due to this background. Therefore the effectiveness of the reaction, particularly for the higher energy levels, will be limited, even if the energy resolution is improved. A further example of this problem is shown in the spectrum from the $(K^-_{\text{stopped}}, \pi^0)$ reaction discussed below.

In addition, kaon absorption generally leads to $\Sigma$ rather than $\Lambda$ production. This is shown in Table IV which is taken from $K^-_{\text{stopping}}$ reactions in bubble chambers [49]. In this table the $R$ factors are the branching fractions to a particular channel upon $K^-$ capture. The ratio, $R_m/R_p$, is the ratio of captures on neutrons to protons. The ratio, $R_{m}$, is the branching ratio for capture on multi-nucleons in the nucleus.

In another study of $K^-$ capture in C, approximately 5 times as many $\Sigma$s as $\Lambda$s were produced [48]. The increase in QF and $\Sigma$ production for stopping $k^-$ reactions will have implications for the FINUDA experiment [50] at DAΦNE. The FINUDA experiment is discussed in another session.

**Table IV. Branching ratios for hyperon production using stopped $K^-$ [49]**

<table>
<thead>
<tr>
<th>Ratio</th>
<th>H</th>
<th>D</th>
<th>He</th>
<th>C</th>
<th>Ne</th>
</tr>
</thead>
<tbody>
<tr>
<td>$R(\Lambda \pi^0)$</td>
<td>4.9</td>
<td>5.</td>
<td>6.2</td>
<td>4.4</td>
<td>3.4</td>
</tr>
<tr>
<td>$R(\Sigma^+ \pi^-)$</td>
<td>14.9</td>
<td>19.</td>
<td>27.3</td>
<td>27.7</td>
<td>27.7</td>
</tr>
<tr>
<td>$R(\Sigma^- \pi^0)$</td>
<td>34.9</td>
<td>22.</td>
<td>10.9</td>
<td>16.8</td>
<td>20.4</td>
</tr>
<tr>
<td>$R(\Sigma^0 \pi^0)$</td>
<td>21.4</td>
<td>23.</td>
<td>21.2</td>
<td>25.7</td>
<td>27.6</td>
</tr>
<tr>
<td>$R(\Lambda \pi^-)$</td>
<td>9.7</td>
<td>10.</td>
<td>12.6</td>
<td>8.7</td>
<td>6.7</td>
</tr>
<tr>
<td>$R(\Sigma^0 \pi^-)$</td>
<td>7.1</td>
<td>5.</td>
<td>5.9</td>
<td>3.3</td>
<td>2.1</td>
</tr>
<tr>
<td>$R(\Sigma^- \pi^0)$</td>
<td>7.1</td>
<td>5.</td>
<td>5.9</td>
<td>3.3</td>
<td>2.1</td>
</tr>
<tr>
<td>$R_{m}/R_p$</td>
<td>0.31</td>
<td>0.25</td>
<td>0.32</td>
<td>0.18</td>
<td>0.12</td>
</tr>
<tr>
<td>$R_m$</td>
<td>0.01</td>
<td>0.16</td>
<td>0.19</td>
<td>0.23</td>
<td></td>
</tr>
</tbody>
</table>

6.3. A *Hypernuclear Production in the $(K^-_{\text{stopped}}, \pi^0)$ Reaction.* - This reaction will be discussed here in more detail as an example of a stopped reaction in which both strangeness and charge are exchanged. In this case, hypernuclear species can be created, charge symmetric to those studied by in-flight, mesonic reactions.

In general, comparison of the spectra of charge symmetric hypernuclei provides information needed to extract the isospin asymmetry of the fundamental, $\Lambda$-nucleon interaction. This has been studied to some extent in the ground states of s- and p-shell mirror hypernuclear pairs, but aside from binding energies, few comparative data are
available[5]. However in addition to charge asymmetry in the fundamental Λ-N interaction, first order Coulomb effects can lead to energy differences between isospin symmetric hypernuclei, in part because the addition of a Λ to the nuclear core changes its radius[51]. Therefore a careful study of the spectra of several charge symmetric pairs is needed to extract both the Coulomb and charge asymmetry effects for the excited as well as the hypernuclear ground states. Indeed, if the Coulomb energy contributions are understood, it should be possible to extract a hypernuclear radius from these data.

In the example to be discussed, stopped kaons are used because essentially all stopped K− interact with the nuclear target, and a stopping beam allows thick targets to be used without degrading the energy resolution. This occurs because the incident momentum is almost always close to zero, and the outgoing particles are photons from π0 decay. Increased production yield is important as the detection efficiency for a π0 is greatly reduced compared to that of charged pions.

In this case the π0 were detected by observing the decay photons using a neutral meson spectrometer, NMS. The NMS [52, 53], is a large acceptance photon detector, and measures the total energy of a π0 by detecting the opening angle and energy of the gamma showers originating in its decay. It consists of 2 CsI arrays of 60 crystals each, fronted by a set of BGO converter and wire chamber tracking planes. The CsI crystals provide the photon calorimetry to determine the relative energy difference between the decay photons, while the BGO and wire chambers determine the location of the photon conversion. This later position, combined with the K− stopping position from a target tracking counter, provides the opening angle between the photons. Figure 19 shows one arm of the NMS assembly. The energy of the π0 is then determined from the opening angle of the decay photons.
Applying a cut on the energy sharing parameter, \( \chi = \frac{E_1 - E_2}{E_1 + E_2} \), which determines the energy difference between the photons from the \( \pi^0 \) decay, improves the energy resolution and background rejection at the expense of acceptance. For a \( \chi = 0.5 \) the \( \pi^0 \) energy resolution with all cuts is 2.6 MeV (FWHM). If this cut is reduced to 0.1 the resolution improves to 1.74 MeV. The cut also improves the spectrum by removing tails on the missing mass peak. A Monte Carlo study has shown that these tails are due to in-flight decays and \( \pi^0 \) contamination from other decay channels.

In the experiment, a dispersed kaon beam with a nominal momentum of 690 MeV/c was brought to rest in a set of 4 natural graphite targets after it traversed a wedge-shaped, brass degrader of central thickness \( \sim 141 \) mm. The targets, each 12.7 mm thick, were components of an Active Target Chamber (ATC) system [54], which provided the transverse as well as longitudinal coordinates of the stopped kaons. The wedge-shaped degrader was used to compensate for the beam dispersion \( \sim 1.2 \frac{MeV/c}{cm} \).

The total energy spectrum of the \( \pi^0 \) emitted in this reaction is shown in figure 20. The overall structure of the spectrum is bounded by the NMS solid angle acceptance, pion production kinematics, and subsequent cuts on various quantities. It can be divided into 1) a low energy region, 2) a region of \( \Sigma \) production, 3) a quasi-free \( \Lambda \) production region, 3) a \( \Lambda \) \( \Lambda \) bound state region, and 4) a kinematically un-allowed region having pion energies above those for \( K^- \) capture at rest forming the hypernuclear ground state. The global shape of the pion energy spectrum is determined by the solid angle acceptance of the NMS.

A predominant channel of hyperon production in the stopped kaon-nucleus interaction is the \( \Sigma \) channel. The two possible ways of producing \( \pi^0 \)'s in association with \( \Sigma \) production are: 1) \( \pi^0 \)'s emitted in association with \( \Sigma^0 \)'s production which follows 3-body kinematics; and 2) \( \pi^0 \)'s produced from \( \Sigma^+ \) decay which can be mono-energetic when the \( \Sigma^+ \) decays at rest. In the later case, a fraction of the charged \( \Sigma \), produced with some recoil momentum, can be brought to rest in the target before they decay. Thus the expected \( \pi^0 \) peak at \( \sim \)
232 MeV, associated with a 2-body $\Sigma^+$ decay at rest is observed in the data. This peak position aids in verifying the energy calibration of the spectrum. The broad structure of the peak is due to $\Sigma^+$ decay in flight.

Background $\pi^0$s can occur through quasi-free $\Lambda$ production or from $\Lambda$ decay. However unless significantly boosted, pions from $\Lambda$ decay have energies either below or near the low-energy cut-off of the NMS acceptance. However $\pi^0$s from quasi-free $\Lambda$ production dominate most of the spectrum.

It is also possible to produce pions with energies above those from the hypernuclear ground state. For example, in-flight $K^-$-nucleus reactions and kaon decays can produce pions in this energy region, although the Lorentz boost requires that energetic pions are forward directed, and therefore their detection is suppressed by the NMS acceptance. However it is also possible that the $K^-$ may be captured on correlated nucleons within the target nucleus, and it has been argued that $\sim 19$ % of all kaon captures occur with multi-nucleons participating in the process [55]. As an example;

\[
(9a) \quad K^- (pp) \rightarrow \Sigma^+ + n; \\
\Sigma^+ \rightarrow p + \pi^0.
\]

The interest of this experiment is to investigate the case when the $\Lambda$ hyperon remains bound to the $^{11}$B nuclear core forming a $^{12}_\Lambda$ hypernucleus. This is an isospin “mirror” hypernucleus to $^{12}_\Lambda$C, which has been previously studied using strangeness exchange and associated production reactions.

The region of the spectrum containing the hypernuclear spectra is re-plotted in figure 21. This region is fit by minimizing a function containing two Gaussian peaks added to a linearly increasing background which is superimposed on a constant background of 1.8 counts per Mev. It yields a strength of $13.7 \pm 4.0$ counts for the ground state at $308.2 \pm 0.2$ MeV, and $17.5 \pm 4.5$ counts for the $(P_3/2, \Lambda_F)$ at 298.0 $\pm$ 0.5 MeV. The energy resolution extracted from the fit was $2.2$ MeV (FWHM), which is larger than expected, given the measured calorimetry and shower positional resolution. This is attributed to the problems associated in maintaining calibrations over the long period of data acquisition.

The binding energy of the $\Lambda$ in the $^{12}_\Lambda$B nucleus is $B_\Lambda = 11.2 \pm 1.0$ MeV, in agreement with the currently accepted value of $11.37 \pm 0.06$ MeV [13].

The background in hypernuclear production using the $(K^-_{\text{stopped}}, \pi^0)$ reaction is larger than for in-flight reactions. The $K^-_{\text{stopped}}$ reaction is not as selective and the actual momentum of the kaon is only assumed, but not measured, to be zero. In addition, pion momenta from both production and decay products occur in the same spectrum region, and multi-nucleon captures also produce unavoidable backgrounds.

The hypernuclear ground state formation probability is $(0.28 \pm 0.08) \times 10^{-3}$ and that for the p-shell of $(0.35 \pm 0.09) \times 10^{-3}$. This is compared in Table V to theoretical and experimental values for the $(K^-_{\text{stopped}}, \pi^-)$ reaction, which should occur twice as often based on isospin conservation. The quoted errors are statistical, but because of the
Fig. 21. – The spectrum in the region of the bound hypernucleus $\Lambda^2$B [54]

difficulty in extracting the yield from the background, the additional systematic error is somewhat larger for the p-shell (about 15%), The formation probability to the ground state is lower than the previous experimental value for $\Lambda^2$C formation after correction for isospin, but remains higher than the theoretical calculations for the ground state [47, 43, 45, 44].

**Table V. – Hypernuclear Formation Probabilities after Stopped $K^-$ Capture**

<table>
<thead>
<tr>
<th>Reference</th>
<th>$\Lambda S_{1/2}$</th>
<th>$\Lambda (P_{3/2} + P_{1/2})$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\Lambda^2$C Theory [43]</td>
<td>$0.33 \times 10^{-3}$</td>
<td>$0.96 \times 10^{-3}$</td>
</tr>
<tr>
<td>$\Lambda^2$C Theory [45]</td>
<td>$0.12 \times 10^{-3}$</td>
<td>$0.59 \times 10^{-3}$</td>
</tr>
<tr>
<td>$\Lambda^2$C Theory [47]</td>
<td>$0.231 \times 10^{-3}$</td>
<td>-</td>
</tr>
<tr>
<td>$\Lambda^2$C Exp [48]</td>
<td>$(0.98 \pm 0.12) \times 10^{-3}$</td>
<td>$(2.3 \pm 0.3) \times 10^{-3}$</td>
</tr>
<tr>
<td>$\Lambda^2$B Theory [47]</td>
<td>$0.119 \times 10^{-3}$</td>
<td>$0.59 \times 10^{-3}$</td>
</tr>
<tr>
<td>$\Lambda^2$B Exp [54]</td>
<td>$(0.28 \pm 0.08) \times 10^{-3}$</td>
<td>$(0.35 \pm 0.09) \times 10^{-3}$</td>
</tr>
</tbody>
</table>

* Multiply by 2 to Compare this Result to $\Lambda^2$C Production

6’4. The ($\pi^+, K^+$) Reaction. – This reaction was first explored at the AGS in a series of investigations providing spectra across a wide range of hypernuclei, with typical resolutions of 3-4 MeV [56, 57]. The reaction was then developed in detail at KEK with a dedicated beamline using a high resolution spectrometer, SKS, specifically built to detect the reaction kaons. Using this spectrometer, resolutions improved to about 2 MeV [58].

The ($\pi^+, K^+$) reaction has so far been the most productive of all hypernuclear reactions. While the cross section is lower than ($K^-, \pi$) reactions, the higher intensities of
pion beams can more than compensate for this deficiency. More importantly, however, the selectivity of the reaction, and the penetrability of the projectiles allow the excitation of selected, deeply lying states. Thus the excited structure is sufficiently separated for nuclei near closed shells, and has sufficient strength, to be identified. The reaction provides a textbook example of the single-particle, shell structure of a nucleus. In figure 22 the particle hole structure of the $^{139}\Lambda$La hypernucleus [59] is clearly evident. Taking the positions of the major shells as observed in this (and some other) reactions, a set of curves of the shell energies vs $A^{2/3}$ provides a classic confirmation [60] that in good approximation, the $\Lambda$ can be considered as an indistinguishable, distinguishable particle when embedded in a nucleus, figure 23.

Counter experiments using magnetic spectrometers have generally provided hypernuclear spectra having energy resolutions $\geq 2$ MeV. This is due to the intrinsic resolutions of secondary mesonic beamlines, and the target thicknesses required to obtain sufficient counting rates. However, a study did achieve a spectrum resolution of approximately 1.5 MeV for the $\pi^+ D(C, K^+)^{12}C$ reaction, using a thin target and devoting substantial time to data collection [59]. This work demonstrated the importance of good resolution, as significantly new information, showing some of the fine structure in the $^{12}_\Lambda C$ spectrum, was obtained 24.

Fig. 22. – The hypernuclear spectrum of $^{139}_\Lambda$La showing the major $\Lambda$ shell structure. [59]
Fig. 23. – The energy levels of the major hypernuclear shells as a function of the $A^{2/3}$ of the hypernucleus. The curve is a calculation using one effective, density dependent potential for the $\Lambda$ nucleus interaction. [60]

Fig. 24. – The $^{\Lambda}n^4\Lambda C$ hypernuclear spectrum with 1.5 MeV energy resolution showing positions of some core excited states. [59]
6.5. The $\gamma, K^+$ Reaction. - Although, specific hypernuclear states can be located within a few keV by detecting de-excitation gammas [61] in coincidence with a hypernuclear production reaction, such experiments become more difficult in heavier systems due to the number of transitions which must be unambiguously assigned in an unknown spectrum. It should also be noted, that resolutions of a few hundred keV are sufficient for many studies, since reaction selectivity and angular dependence potentially allows extraction of the spectroscopic factors to specific states [62]. A reaction also provides a full spectrum of states which can be clearly identified with a specific hypernucleus. Indeed the excitation strength of the spectrum is of interest, as it directly relates to the model that the reaction proceeds through the interaction of the incident projectile with an identifiable nucleon within the nuclear medium. Thus as an example apropos to the experiment reported here, if the theoretical spectrum does not reproduce the experimental one, it is possible that propagator re-normalization within the medium could be significant [63], requiring a modification of the single particle picture of the reaction.

Electroproduction of hypernuclei is illustrated in figure 6. Electroproduction traditionally has been used for precision studies of nuclear structure, as the exchange of a colorless photon can be accurately described by a first order perturbation calculation. In addition, electron beams have excellent spatial and energy resolutions. Previously, electron accelerators had poor duty factors, significantly impairing high singles rate, coincidence experiments. However, modern, continuous beam accelerators have now overcome this limitation, and although the cross section for kaon electroproduction is some 2 orders of magnitude smaller than hadronic reactions, this can be compensated by increased beam intensity. Targets can be physically small and thin (10-50 mg/cm²), allowing studies of almost any isotope. The potential result for $(e, e'K^+)$ experiments, is an energy resolution of a few hundred keV with reasonable counting rates up to at least medium weight hypernuclei [64].

Finally, the $(e, e'K^+)$ reaction, because of the absorption of the spin 1 virtual photon, has high spin-flip probability even at forward angles. In addition the momentum transfer is high, approximately 300 MeV/c at zero degrees for 1500 MeV incident photons, so the resulting reaction is expected to predominantly excite spin-flip transitions to spin-stretched states of unnatural parity [66]. These states are not strongly excited in mesonic production, and the electromagnetic process acts on a proton rather than a neutron creating proton-hole Λ-states, charge symmetric to those previously studied with meson beams. Precision experiments, comparing mirror hypernuclei, are needed to extract the charge asymmetry in the ΛN potential.

In electroproduction, the energy and momentum of the virtual photon are defined by $\omega = E(e) - E(e')$ and $\vec{q} = \vec{p}(e) - \vec{p}(e')$, respectively. The four-momentum transfer of the electron is then given by $Q^2 = q^2 - \omega^2$. Since the elementary cross section, and particularly the nuclear form factor, fall rapidly with increasing $Q^2$, experiments must be done within a small angular range around the direction of the virtual photon. In addition as discussed below, the virtual photon flux is maximized for an electron scattering angle near zero degrees [67, 68]. Thus the experimental geometry requires two spectrometer arms, one to detect the scattered electron and one to detect the kaon, both placed at
extremely forward angles.

To a good approximation, the electroproduction cross section can be expressed [65] by;
\[
\frac{\partial^3 \sigma}{\partial E \partial \Omega_k \partial \Omega_k} = \Gamma \left[ \frac{\partial \sigma_T}{\partial \Omega_k} + \epsilon \frac{\partial \sigma_L}{\partial \Omega_k} + \epsilon \frac{\partial \sigma_R}{\partial \Omega_k}\cos(2\phi) + \cos(\phi_k) \sqrt{2\epsilon(1+\epsilon)} \frac{\partial \sigma_L}{\partial \Omega_k} \right].
\]

Here \( \phi \) is the out of plane angle, and the factor, \( \Gamma \), is the virtual flux factor evaluated with electron kinematics in the lab frame. It has the form;
\[
\Gamma = \frac{\alpha}{2\pi^2 Q^2} \frac{E_\gamma - E'_e}{1 - \epsilon E'_e}.
\]

In the above equation, \( \epsilon \) is the polarization factor;
\[
\epsilon = \left[ 1 + \frac{2k^2}{Q^2} \tan^2(\Theta_e/2) \right]^{-1}.
\]

For those virtual photons almost on the mass shell, \( Q^2 = p_\gamma^2 - E'_\gamma \to 0 \). The label on each of the cross section expressions \( (T, L, P, \text{ and } I) \) represent transverse, longitudinal, polarization, and interference terms. For real photons of course, \( Q^2 \to 0 \), so only the transverse cross section is non-vanishing, and for the experimental geometry used here, the cross section is completely dominated by the transverse component. The electroproduction cross section may also be replaced, to good approximation, by the photoproduction value.

Experimentally, \( \Gamma \) is integrated over the angular and momentum acceptances of the electron spectrometer. In order to maximize the elementary, \( p(\gamma, k^+)\Lambda \) reaction, the photon energy is chosen to be about 1.5 GeV. In addition, to keep strangeness production limited essentially to kaons and \( \Lambda_s \), the energy, \( E'_e \), of the incident electron is set to be \( \approx 1.8 \text{ GeV} \). In this way, backgrounds from unwanted reactions are reduced. This also allows a physically small, low-momentum electron spectrometer to be employed, as the scattered electron energy, \( E'_e \), is about 0.3 GeV.

Figure 25 shows the calculated virtual photon flux factor in units of photons per electron per MeV-sr for the chosen kinematics. This flux factor peaks near zero degrees and falls rapidly as the scattering angle increases. With electrons detected at zero degrees, a large percentage of the scattered electrons are captured by even a small solid angle, increasing the coincidence probability between these electrons and the reaction kaons of interest. In addition, because of the small beam spot, \( \approx 100 \mu m \), the \( \approx 0^\circ \) electron scattering angle, and the small momentum value of the scattered electron, it is sufficient to only measure the electron position on the spectrometer focal plane to ensure excellent energy resolution. However, the disadvantage of this geometry is a high electron background rate from target bremsstrahlung, which ultimately limits the usable beam luminosity.

One can choose to place the electron spectrometer at a small but finite angle [69], or at zero degrees [70]. A spectrometer at finite angle requires high luminosity and large acceptance due to the lower value of the virtual photon flux. Here the zero degree
Fig. 25. – The virtual photon flux factor as a function of the electron scattering angle. [68]

detection geometry is developed in more detail below. Once the choice of the incident and scattered electron momenta are fixed, the reaction kaon momenta are determined by the kaon reaction angle. In this experiment, the chosen kinematics produced a kaon momenta of $\approx 1.2 GeV/c$, providing a 3-momentum transfer of $\approx 300 MeV/c$ to the recoiling $\Lambda$. The kaon momentum provides reasonable kaon lifetime, and allows $\pi/K$ discrimination using threshold aerogel Cerenkov detectors coupled with time of flight. Figure 26 shows a schematic view of the experimental layout.

The beam has a bunch width of 1.67 ps with a bunch separation of 2 ns. While the absolute value of the energy of the incident electrons was unimportant (although for kinematic reasons to be discussed below, it did need to be determined), it was extremely important to precisely maintain whatever this energy was over the several weeks of the
experiment. Thus the beam momentum is locked by a "Fast Feedback Energy Lock System" installed in an arc of the accelerator. This system measured, at a repetition rate of 1 kHz, the beam parameters at the entrance, the position of maximum momentum dispersion, and the exit of the arc, to extract an energy correction factor. This correction was then applied to the last cavity of the accelerator, maintaining a constant beam energy. The feedback lock controlled the total energy of the primary electron beam to a \( \delta p/p \leq 10^{-4} \). The intrinsic energy spread in the beam was controlled by a tune of the injector. A similar lock system maintained the beam position on target within 100 \( \mu \)m. Although the intrinsic spot size was tuned to be \( \leq 100 \mu \)m, the beam was de-focused to \( 4 \times 4 \) mm\(^2\) by a fast raster when incident on the CH\(_2\) target, to reduce beam heating.

Beam intensities were set to produce an acceptable signal to accidental ratio, which for the C target was approximately 0.6 \( \mu \)A, or an experimental luminosity of approximately \( 4 \times 10^{33} \) cm\(^{-2}\)s\(^{-1}\). To protect the CH\(_2\) target, the beam current was kept below 1.5 \( \mu \)A.

In order to detect both scattered electrons and positively charged kaons near zero degrees, a "C" magnetic dipole (splitter) was used. The target was positioned at the upstream side of the effective field boundary of this magnet. The splitter respectively deflected electrons scattered at approximately 0 degrees and kaons at approximately 2 degrees by 33 and 16 degrees, respectively.

A short orbit spectrometer (SOS) is one of two existing magnetic spectrometers in Hall C at JLab, and as it has a flight path of \( \approx 10 \)m, this spectrometer is particularly useful for the detection of particles with short half-lives. However, it has low dispersion and large momentum acceptance, and these characteristics are not well matched to the present experimental geometry. Still the SOS was mounted at the Hall C pivot, and had the sophisticated particle identification (PID) package\(^{[72]}\) required to identify kaons within the large background of pions and positrons, so it was chosen as the kaon spectrometer for this first \((e,e'K^+)\) experiment. It was expected that the overall resolution would be dominated by this spectrometer\(^{[76]}\).

The solid angle acceptance of the splitter/SOS spectrometer system was approximately 5 msr, covering a range of scattering angles from 0 to 4 deg. The error in the reconstructed scattering angle was about 13 msr (FWHM), and was dominated by the horizontal angular error measurement. This contributed about 200 keV to the missing mass resolution when the recoil mass was \( \geq 6 \). The central momentum of the SOS was set to 1.2 GeV/c. The acceptance was \( \approx 46\% \), but only the central \( \pm 15\% \) was useful. This acceptance was nearly flat within the missing mass range of interest.

The standard SOS detector package was used. It consisted of: 1) two sets of tracking chambers separated by 0.5 meters; 2) four scintillation hodoscope planes; 3) one aerogel Cerenkov (AC) counter with an optical index of 1.03; 4) one Lucite total internally reflective Cerenkov (LC) counter with index 1.49; 5) one gas Cerenkov (GC) detector; and 6) 3 layers of lead-glass shower counters. The tracking detectors were used to determine the position and angle of the particle on the focal plane and by projection, its scattering angle from the target. The scintillator hodoscopes were used to localize tracks in the wire chambers, and to obtain timing and time of flight (TOF) information for PID. The aerogel Cerenkov detector was used to veto pions and positrons, and the Lucite counter
was used to remove protons by tagging high-beta particles. The gas Cerenkov detector and the lead-glass shower counters were used to remove positrons.

It was expected that the numbers of positrons, pions and protons in the SOS would be very much larger than the number of kaons. Indeed the measured flux per second of positrons, pions, protons, and kaons from the C target, was $10^5$, $1.4 \times 10^3$, 140, and 0.4, respectively. Therefore excellent particle identification was required, not only in the analysis, but also in the hardware trigger.

The standard SOS detector package was used to identify kaons as discussed above. The large flux of positrons was due to the acceptance of scattering angles down to zero degrees, where positrons from Dalitz pairs, created in the target, were observed. Positrons were easily identified and could have been removed in the trigger by the lead-glass shower counter, but detection of the Dalitz pairs provided a useful confirmation of the experimental resolution.

The coincident time resolution between an electron and a kaon was 230 (FWHM) ps, after pulse height and path length corrections were applied. Coupled with a measure of the kaon time-of-flight, the system time resolution was sufficient to identify the real and accidental coincidence peaks, and the true kaon coincident events were selected by a two-dimensional cut on a real coincidence window of 2 ns as shown in figure 27. Events selected from an average of eight nearby accidental coincidence windows were used to obtain the shape and magnitude of the accidental background spectrum.

The scattered electrons were detected in a split-pole, magnetic spectrometer[73], which was well matched to the geometrical kinematics and acceptances. The spectrom-
eter coupled with the phase space of the incident beam, had the capability of obtaining $5 \times 10^{-4}$ resolution (FWHM $\delta p/p$). The central momentum was chosen to be 300 MeV/c with a momentum acceptance of $\approx 120$ MeV/c. The solid angle acceptance of the combined splitter/split-pole system was about 9 msr, which effectively tagged about 35% of the virtual photon flux. This provided a spread in the flux momentum of $\approx 120$ MeV/c, centered around $\approx 1500$ MeV/c. In summary, the geometry of the electron arm was possible because of the excellent phase space of the incident electron beam, the thin targets which limited multiple scattering, and the extremely forward peaking of the virtual photon flux factor.

However, target bremsstrahlung also peaks at zero degrees and large numbers of scattered electrons are expected to enter the split-pole spectrometer[68]. In fact the experimental luminosity was set by accepting a total rate of $\approx 2 \times 10^8$ on the instrumented portion of the focal plane. To operate at this rate, the detection system required that only the focal plane position of a scattered electron is needed to obtain the required resolution, since more detailed tracking would have been very inefficient. The focal plane detector[74] was composed of 10, one-dimensional silicon strips segments (SSD), each having 144 strips with a pitch of 0.5mm. These segments were placed approximately perpendicular to the electrons, which were incident at $\approx 47^\circ$ on a 72 cm length of the the focal plane. The singles rate per strip was $\approx 10^6$ s$^{-1}$.

A set of 8 scintillation strip counters in a hodoscope arrangement were positioned directly behind each of the SSD segments. These strips were 1cm wide, 6cm long and 0.4 cm thick, viewed at one end through a light guide by a 3469 Hamamatsu photomultiplier. Rates per scintillator were found to be $\leq 1.5 \times 10^6$ hz, and no change in time resolution was observed up to rates $\leq 2.5 \times 10^6$ hz. The SSD provided the position of an electron event to within 500 $\mu$m and the scintillation hodoscope provided event timing to 250 ps ($\sigma$).

Since the entire beamline/spectrometer system was under vacuum, multiple scattering in the air and vacuum windows occurred only at the exit of the spectrometers. Vacuum windows were located immediately before the first tracking detectors so that effects on the measured track-position were minimized. Table 1 lists the expected contributions to the energy resolution. As discussed above, the contribution from the SOS spectrometer was expected to dominate.

Analysis of the experiment required knowledge of the magnetic transport coefficients of the spectrometers. Although the coefficients for the SOS transport had been previously established in several experiments, the addition of the splitting magnet to the system required that they be re-determined. For example, the angular acceptance of the SOS, normally 7 msr for a point target, was reduced to 5 msr by the splitter.

The reconstructed missing mass of a hypernuclear state is a function of the beam energy, the momenta of the scattered electron and kaon, and the scattering angles. In a two-dimensional space defined by the electron and kaon momenta, the recoil missing mass is obtained by a projection of the events onto a locus line. Using an incorrect value of the beam energy or central momentum value for either spectrometer arm, results in an incorrect position and slope of the locus line, and therefore an incorrect kinematic
position and width for various missing masses. Thus one not only needs the relative values of scattering angles and the fractional change of the scattered momentum with respect to central momenta, but also the absolute values of the beam energy and central momenta of the spectrometers, or the absolute values of the scattering angles of the coincident particles.

The simulation of charged particle trajectories through the system of magnets and detectors used the program, RAYTRACE[75], and the coefficients of the RAYTRACE code were determined by adjusting the splitter contribution so that the calculated multi-dimensional, phase-space distributions from a point beam matched those measured when the entrance angles and positions of reaction protons and pions from the target were restricted by a set of appropriately positioned holes in a tungsten plate (sieve slit) located between the splitting magnet and the SOS[76]. Optimization of the SOS coefficients used a χ² minimization process defined by the difference between the simulated and observed experimental patterns.

The binding energy spectrum of the $^{12}$B hypernucleus is shown in figure 28 with the backgrounds. Two prominent peaks are obvious in the spectrum. This is similar to that predicted by Motoba, et al[77] and by Millener[78]. Reference [77] calculates the excitation strengths in DWIA for the photoproduction process of a kaon at an angle of 3° by a 1.3 GeV photon. The curve in the figure is generated by superimposing Gaussian peaks, having a FWHM of 750 keV below and 5 MeV above 15 MeV excitation energy, on a polynomial fit to the averaged accidental background. The positions of the states are taken from ref [78], as this latter spectrum was obtained from an effective p shell Λ-nucleus interaction previously matched to $(\pi^+, K^+)$ data. The reaction strengths [77] for the low-lying states of $^{12}$ΛB are shown in figure 29. The theoretical curve in figure 28 is directly overlayed on (not fitted to) the data.

The major excitations are in good statistical agreement with theory both in position and strength. However, the core excited states, predicted to lie between the major shell excitations, seem more weakly excited than predicted, but statistics are not sufficient to discuss this region of the spectrum in detail. The differential cross section can be calculated as if it were photoproduction, by assuming the virtual photons are massless.

### Table VI. Contributions to the System Energy Resolution

<table>
<thead>
<tr>
<th>Source</th>
<th>Contribution</th>
<th>Resolution (keV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Beam Energy</td>
<td>$10^{-3}$</td>
<td>$\leq 180$</td>
</tr>
<tr>
<td>SOS momentum</td>
<td>$5.5 \times 10^{-3}$</td>
<td>$\approx 660$</td>
</tr>
<tr>
<td>Split-pole</td>
<td>$5 \times 10^{-4}$</td>
<td>150</td>
</tr>
<tr>
<td>Kaon Scattering Angle $(^{12}C)$</td>
<td>13 mr</td>
<td>$\approx 200$</td>
</tr>
<tr>
<td>Target Energy Loss $(^{12}C)$</td>
<td>1.7 keV/mg/cm²</td>
<td>38</td>
</tr>
<tr>
<td>Total</td>
<td></td>
<td>$\approx 757$</td>
</tr>
</tbody>
</table>
This averages the elementary (γ,K) reaction at 1500 MeV over the ≈ 100 MeV spread of virtual gamma energies. The weighted average of the cross section measurements for the ground state doublet at the two incident beam energies is, 140 ± 17(stat) ± 18(sys) nb/sr. The value is consistent with the theoretical prediction[66] of ≈ 138 nb/sr.

The 2− state of the ground state doublet lies at approximately 150 keV excitation energy, and is expected to be dominate. The resolution (and statistics) is insufficient to identify this splitting. The 3+ p-shell state is also predicted to dominate the spectrum in the 10 MeV excitation region. While theory indicates one dominate peak at this excitation energy; this peak is created by the predicted overlap of the 2+ and 3+ states in the the p-shell region[78]. These two states are not as degenerate if other [77] effective parameters are used. The results demonstrate sensitivity to the effective interaction and the DWIA transition amplitudes.

To confirm the cross section normalization, the quasi-free component of the experimental spectrum was extracted, and the yield corrected for acceptance and momentum transfer [80]. One obtained 4.2 interacting protons, in agreement with previous measurements [81, 82].

6.5.1. Final State Interactions. Photoproduction of QF hyperons in the impulse approximation can be written in the form;
Fig. 29. – A schematic representation of the reaction strength for the low lying states of $^{12}$B [77]

\[
\frac{d\sigma}{d\Omega d\Omega_{\text{N}} d\Omega_{\text{Y}}} = (\text{kinematic factor}) \sum |T_{\gamma p} M_{fi}| \\
\delta(\omega + M_T - E_N - E_Y - E_K) \delta(\vec{q} - \vec{p}_N - \vec{p}_Y - \vec{p}_K);
\]

where the energy (momentum) definitions are contained in figure 30. The transition matrix contains the elementary amplitude and the spectroscopic factor for the nuclei involved. This is written in a plane wave approximation for the reaction, $\gamma d \rightarrow K^+YN$. Using the figure 30:

\[
T_{FI} = \langle e^{i\vec{k}_F \cdot \vec{R}} \phi(\vec{r}) | V(\vec{r}_p) | e^{i\vec{k}_F \cdot \vec{R}} \psi_d(\vec{r}) \rangle
\]

In the $t\rho$ approximation, $T_{FI}$ is written as the product of the elementary amplitude, $t_{\gamma p}$ times the nuclear matrix element:

\[
M_{fi} = \int d^3r \phi^{\ast} e^{-i\vec{q} \cdot \vec{r}} \psi_d;
\]

where $\vec{q} = (\vec{k}_F - \vec{k}_F)/2$. This approximation is useful when both $\vec{q}$ and the relative energy between the hyperon and nucleon are small. Experimental data [83] for the $\gamma d \rightarrow K^+\Lambda N$ reaction are shown in figure 31. Here one notes that photoproduction of $\Lambda$s involves a spin flip of the $n \rightarrow \Lambda$ transition, so the resultant final state interaction will be in a
spin singlet state. The strength of the FSI enhancement is then proportional to the s-wave, spin singlet scattering length and effective range. For the present data, the energy resolution and statistics are insufficient to extract these values. Substantially better resolution is possible, and the experiment should be repeated.

6’6. The \((N; N', K^+)\) Reaction. This reaction is similar to electroproduction, with virtual mesons being the exchanged quanta. These are produced in \((N, N')\) scattering, and both pion and kaon exchange are needed to explain the elementary strangeness production process near threshold, figure 32. However, the dominant mechanism seems to be pion exchange, with kaon exchange required to reproduce the structure which is observed at the \(\Sigma\) threshold.

That the \((N; N', K^+)\) reaction has not been seriously applied to hypernuclear pro-

Fig. 31. – The experimental spectrum showing the \(\Lambda\) FSI predictions for various spin singlet scattering lengths. [83]
duction, is due to the lack of appropriate facilities, although several experiments have produced light, s-shell hypernuclei [84, 85]. The momentum transfer is large[86], but is comparable to $(\pi, K)$ when the exchanged pion is nearly on the mass shell. However, for incident momenta above the $\Lambda$ threshold, the nucleons and kaon have momenta of several GeV/c, requiring large, specially designed spectrometers. As with the $(e,e'K^+)$ reaction discussed above, the dual nucleon and kaon spectrometers would both need to both cover the very forward laboratory angles. Resolutions would probably be poor due to the high momentum of the detected particles. Although cross sections [88], figure 33, are a few tens of nb/MeV·sr, proton beam intensities are sufficient for such experiments. Of more importance, is the expected background, especially QF production at the momentum transfers of the reaction, and the selectivity of the states produced.

67. The $(p, K^+)$ Reaction. – The $(p, K^+)$ reaction [89] results in the production of hypernuclei by the direct insertion of a $\Lambda$ without the creation of a hole state. Cross sections are estimated as shown in figure 34. The important lesson here is the effect of $NN$ correlations which can lead to a factors of $10^2-3$ increase in the cross sections. Finally given the low cross section, and difficulty in interpretation of large momentum

Fig. 33. – The kaon spectrum at 10 deg from the reaction $p(p, K^+)YN$ for incident protons of 2.3 GeV energy. [88]
transfer data, this production mechanism is unlikely to be used.

6.8. The Weak Production of Strangeness in NN Collisions. – Although experimentally less accessible than the weak decay of strange systems, non-leptonic weak production of strangeness provides greater control over the reaction channels, and thus greater ability to resolve the reaction mechanism. The integrated cross section for the reaction $pN \rightarrow p\Lambda$ is shown in figure 35. Note that the peak at about 1150 MeV/c is due to the inclusion of $\Sigma N$ coupling through $\rho$ exchange. Pion exchange dominates at other momenta. The production is dominated by parity conserving transitions in the ratio 3/1. Thus an attempt to use parity violation to reduce background will also reduce the signal.

The very low cross section[90] makes an experiment extremely difficult, although at least one group is considering such a measurement [91]. A specific design to observe the reaction $pn \rightarrow p\Lambda$ would use a deuteron beam to provide the neutron, with the proton acting as a spectator. The maximum weak production cross section would occur at about 2.25 MeV/c incident momentum. In addition to detecting the outgoing protons which
define the kinematics, the decay of the recoiling $\Lambda$ some 5 cm downstream of the target would be used to reduce background. The experiment would need to have a signal to noise discrimination level $> 10^{-13}$.

7. – The Production of $S = -2$ Nuclear Systems

Strangeness -2 could be introduced directly into a nucleus by capture of a $\Xi$. Ranging $\Xi^{-}$ capture into atomic orbitals, and eventually interact with the nuclear core before they decay. Unfortunately, the energy transfer to the hyperons in the strong interaction $\Xi N \rightarrow YY'$ is 29 MeV, leading with large probability to QF production of at least one of the hyperons. In addition, direct production of $\Xi$ using the $(K^{-}, K^+)$ reaction requires a two step process, with the transfer of a strange quark from the $K^-$ and the associated production of an $s\bar{s}$ pair. In general this produces $\Xi$s recoiling with momenta that are too high for most of them to range before they decay.

However, there has been an intensive search for the dibaryon composed of two $u$, $d$, and $s$ quarks, whose SU(3) symmetry energy is calculated to produce a deeply bound system, the H particle [92, 93, 94]. In the hadronic limit, the H might be considered as a bound (or low lying resonant) composition of $\Lambda\Lambda$, $\Sigma\Sigma$, and $\Xi N$ channels. Still as usually discussed, the "H" is considered as a 6 quark state. No evidence for an H particle has been found in experiments looking at both production and decay channels [93] across the $\Lambda\Lambda$ threshold, 2.23 GeV.

As will be discussed below, double $\Lambda$ hypernuclei have conclusively been observed through the sequential pion decays [95, 96, 97, 98, 99] of their s-shell $\Lambda$s. The impact of this on the existence of the H is not completely understood, but in principle a deeply bound H could form from the close proximity of the hypernuclear $\Lambda$s, destroying the hypernucleus while creating an H. Arguments about barrier penetrations notwithstanding, the comparatively long lifetime of double $\Lambda$ systems seems to preclude the existence of a deeply bound H dibaryon. The non-existence of an H does not reflect on the possibility of a hadronic state held together by meson exchange.
71. The Double $\Lambda$ systems formed in ($K^-, K^+$) Reaction. - There is very little experimental information on multiply strange systems. Only four $\Lambda\Lambda$ hypernuclei are reported, as single events in five experiments, and one of the observations could not uniquely identify the hypernucleus, Table VII. The $^6_{\Lambda\Lambda}He$ was seen in two different experiments.

The binding energies of these events are not consistent, and the interpretation of the counter experiment involves a rather "unusual" decay mode for its explanation. The internal consistency of the data is improved by neglecting the earlier, and retaining the more recent, $^6_{\Lambda\Lambda}He$ event. Theoretically the NSC97 Nijmegen model then reproduces the experimental binding energy of this reduced data set, although other Nijmegen models cover a range of possibilities. If the reduced data set is accepted, the $\Delta B_{\Lambda\Lambda}$ is changed from approximately 4.6 to 1 MeV; i.e., from a strong binding to a weak one. The quantity, $\Delta B_{\Lambda\Lambda}$ is defined by the equations;

\[
\begin{align*}
\text{Mass}(\Lambda\Lambda A) &= \text{Mass}(A - 2) + 2 \text{Mass}(\Lambda) - B_{\Lambda\Lambda} \\
\Delta B_{\Lambda\Lambda} &= 2 B_A(A - 1) - B_{\Lambda\Lambda}
\end{align*}
\]

Thus $\Delta B_{\Lambda\Lambda}$ represents the additional binding energy in the system which comes from the mutual interaction of the two $\Lambda$s in the nucleus. The weak $\Delta B_{\Lambda\Lambda}$ value confirms the fact that an $H$, at least as a hadronic state, should not exist. In fact it is questionable if even $^4_{\Lambda\Lambda}H$ is bound [102]. Predicted binding energies of double $\Lambda$ systems are given in Table VIII. This table is based on theoretical predictions, and assumes that $B_{\Lambda\Lambda}$ is an approximate constant 1 MeV for $A > 5$.

<table>
<thead>
<tr>
<th>Hypernucleus</th>
<th>Detection</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^\Lambda_{\Lambda}Be$</td>
<td>Emulsion</td>
<td>[95]</td>
</tr>
<tr>
<td>$^6_{\Lambda\Lambda}He$</td>
<td>Emulsion</td>
<td>[96]</td>
</tr>
<tr>
<td>$^{\Lambda}<em>{\Lambda}Be$ or $^\Lambda</em>{\Lambda}B$</td>
<td>Emulsion</td>
<td>[97]</td>
</tr>
<tr>
<td>$^\Lambda_{\Lambda}Be$</td>
<td>Hybrid Emulsion</td>
<td>[98]</td>
</tr>
<tr>
<td>$^4_{\Lambda\Lambda}H$</td>
<td>Counter</td>
<td>[99]</td>
</tr>
</tbody>
</table>

The counter experiment [99] which reported the observation of the $^4_{\Lambda\Lambda}H$ hypernucleus is now discussed in more detail. The experiment used the ($K^-, K^+$) reaction on a $^9Be$ target. Because most of the $\Xi^-$ recoils decay before they range in the target, the reaction is expected to proceed via multi-nucleon interaction in a nucleus, with nucleon(s) emission(s). Probabilities for such processes are given in Table IX. The double $\Lambda$ systems $^6_{\Lambda\Lambda}He$ and $^7_{\Lambda\Lambda}He$ are expected to be most copiously produced [100].

The experiment placed a wide mass cut on the missing mass region of the ($K^-, K^+$) spectrum where double $\Lambda$ hypernuclei are expected to form. Then the momentum of the
Table VIII. - Binding Energies of Double $\Lambda$ Hypernuclei

<table>
<thead>
<tr>
<th>Hypernucleus</th>
<th>$B_{\Lambda}(A-1)$</th>
<th>$B_{\Lambda\Lambda}$</th>
<th>$\Delta B_{\Lambda\Lambda}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{4}_{\Lambda}\Lambda H$</td>
<td>0.13</td>
<td>-0.12</td>
<td>-0.14</td>
</tr>
<tr>
<td>$^{5}_{\Lambda}\Lambda H$</td>
<td>2.04</td>
<td>3.26</td>
<td>0.82</td>
</tr>
<tr>
<td>$^{6}_{\Lambda}\Lambda He$</td>
<td>2.39</td>
<td>3.8</td>
<td>0.98</td>
</tr>
<tr>
<td>$^{6}_{\Lambda}\Lambda He$</td>
<td>3.12</td>
<td>7.25</td>
<td>1.0</td>
</tr>
<tr>
<td>$^{6}_{\Lambda}\Lambda He$</td>
<td>4.18</td>
<td>9.36</td>
<td></td>
</tr>
<tr>
<td>$^{7}_{\Lambda}\Lambda Li$</td>
<td>4.50</td>
<td>10.0</td>
<td></td>
</tr>
<tr>
<td>$^{8}_{\Lambda}\Lambda Li$</td>
<td>5.58</td>
<td>12.16</td>
<td></td>
</tr>
<tr>
<td>$^{9}_{\Lambda}\Lambda Li$</td>
<td>6.80</td>
<td>14.60</td>
<td></td>
</tr>
<tr>
<td>$^{7}_{\Lambda}\Lambda Be$</td>
<td>6.84</td>
<td>14.93</td>
<td>1.25</td>
</tr>
<tr>
<td>$^{10}_{\Lambda}\Lambda Be$</td>
<td>6.71</td>
<td>15.65</td>
<td>2.1</td>
</tr>
</tbody>
</table>

Pions from sequential, two-body mesonic decays of the hypernuclei was observed. For example, the decay sequence was:

$$\gamma Y A \rightarrow \pi + (Y A) \rightarrow \pi + (\pi + A)$$

The pion momenta were determined by a cylindrical magnetic spectrometer surrounding the target, and tracking of the pions in the magnetic field, pulse height, and timing provided particle identification. The two dimensional, coincident pion spectra is shown in figure 36. In this figure the higher of the two pion momenta, $p_h$, is plotted on the vertical axis and the lower, $p_l$ on the horizontal axis. The box size represents the number of events at those momenta.

The figure shows several regions of enhanced counts. The region 1, positioned near pion momenta (114,132) MeV is easily identified as due to the production of a fissioned system of the two single hypernuclei, $^{3}_{\Lambda}H$ and $^{4}_{\Lambda}H$ from their mono-energetic decays. The open question concerns the enhancement at pion momenta near (105,115) MeV, region 2. There are no possible single $\Lambda$ hypernuclear decays which fit the coincidence of these decay momenta. Also while the decay of the $^{4}_{\Lambda\Lambda}H$ system to the ground state of $^{4}_{\Lambda}He$ is close, its decays do not match the observations. Although some states in $^{4}_{\Lambda}Li$ may be partially isomeric, calculations show that the combined decays from all levels in $^{7}_{\Lambda}Li$ populated in decays from the original double hypernucleus cannot explain the spectra as well.

Thus in order to fit the observation, a resonant state [101] was invoked in the $^{4}_{\Lambda}He$ system so that the pion decays proceeded as;

$$^{4}_{\Lambda\Lambda}H \rightarrow \pi + ^{4}_{\Lambda}He* (8.9)$$

$$^{4}_{\Lambda}He* \rightarrow p + ^{3}_{\Lambda}H$$
TABLE IX. ~ Formation Probabilities in the two-step \((K^- , K^+)\) Reaction

<table>
<thead>
<tr>
<th>Hypernucleus</th>
<th>Formation Probability</th>
</tr>
</thead>
<tbody>
<tr>
<td>(^4\Lambda H)</td>
<td>0.02</td>
</tr>
<tr>
<td>(^5\Lambda H)</td>
<td>0.07</td>
</tr>
<tr>
<td>(^6\Lambda He)</td>
<td>1.0</td>
</tr>
<tr>
<td>(^7\Lambda He)</td>
<td>25</td>
</tr>
<tr>
<td>(^8\Lambda He)</td>
<td>26.1</td>
</tr>
</tbody>
</table>

\[ ^3\Lambda H \rightarrow \pi + ^3\Lambda He \]

The resonance assignment was rather arbitrary, and selected to match the observed pion energies. There is no other evidence for such a state, and it certainly is unlikely to find a narrow structure this high in excitation in this hypernucleus. Thus matching the observed pion decay energies using consistent values of \(\Delta B_{\Lambda\Lambda}\) and the known \(B_{\Lambda}\) values of single \(\Lambda\) hypernuclei, presents a difficult puzzle.

However, a reasonable representation of the data can be obtained by superimposing decays from \(^4\Lambda H\) to those from \(^7\Lambda He\). The simulation must carefully include decays to and from all levels, including the continuum. In this representation a \(\Delta B_{\Lambda\Lambda} = 1\) Mev

Fig. 36. – The coincident pion spectrum decays emitted in the reaction \(^9\Lambda Be(\pi^-, \pi^+)\) \(\Lambda\) Reaction
was assumed. The spectra and the normalized simulations are shown on the projected pion axes, \( p_h \) and \( p_t \), in the figure 37. The more narrow structure of the \( p_t \) spectrum is due to the cut-off when \( p_h = p_t \).

Finally, this later explanation requires that \( \Lambda \Lambda \) is bound, although the theoretical prejudice is that it is not [102]. It does, however, support the weak value of \( \Delta B_{\Lambda \Lambda} \), and does not require the postulate of an “unusual” resonance.

The spectra also show some evidence for other double \( \Lambda \) systems in the data, ie region 3, but statistics are too low to draw definitive conclusions. Experimentally much better energy resolution is needed and some type of identification of the decaying system would be helpful in future experiments of this sort.

772. Nuclear Systems containing a \( \Xi^- \). - The existence of \( \Xi \) hypernuclei is not established, although there are several emulsion events [103, 104] which can be interpreted as \( \Xi^- + C \rightarrow \Lambda \Lambda \) \( + \Lambda \) \( + \Lambda \Lambda \) \( + \Lambda \). In one of these events [103], the binding of the intermediate \( \Xi \) state before fission was interpreted as \( B_{\Xi} = 0.54 \) MeV, probably indicating that fission occurred from an atomic level. In the other event [104], the binding energy was indeterminate because the recoiling hypernuclei could have been in an excited state. However, a binding energy as large as 3.7 MeV could have been possible.

The missing mass spectrum in another experiment [105] using the \( nat C(K^-, K^+) \) reaction was observed with very poor resolution. Individual levels could not be ascertained, but from QF analysis, the \( \Xi \)-nucleus well depth was estimated to be \( \approx 15 \) MeV. If this is correct, a bound state should exist in the \( \Xi^- 14 \) \( B \) system with a binding energy of about 6 MeV. The width of the \( \Xi \) states depends on the strong conversion \( \Xi N \rightarrow \Lambda \Lambda \). Still widths have been calculated to be no more that a few MeV, and cross sections and production rates are accessible with present technology. A spectroscopy of \( \Xi \) hypernuclei thus seems possible.
Fig. 38. – The coincident pion spectrum decays emitted in the reaction $^9\text{Be}(K^-, K^+)X$ reaction [113]

8. – Hypernuclei formed in Heavy Ion Collisions

The multiplicity of particles, particularly pions, is large in relativistic heavy ion collisions, and these can produce multiple hyperons which can cool with fragmentation or coalesce into hypernuclei[113]. Cross sections for $\Lambda$ hypernuclei are given in Table X using a coalescence model. The energy dependence of the cross section is shown in figure 38. The coalescence model is based on the probability of finding particles with nearly the same momenta, in the same spatial location. The coalescence factor, $S_{AA}$, is then determined from the overlap of the $\Lambda - A$ wave function with the $\Lambda$ and $A$ source wave functions.

Table X. – Predicted cross sections for the formation of hypernuclei in $\text{Si} + \text{Au}$ collisions at 14.5 GeV/nucleon [4]

<table>
<thead>
<tr>
<th>Hypernucleus</th>
<th>Peripheral Collisions</th>
<th>Central Collisions</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>$\mu$b</td>
<td>mb</td>
</tr>
<tr>
<td>$^4\Lambda$H</td>
<td>0.15</td>
<td>0.34</td>
</tr>
<tr>
<td>$^4\Lambda$He</td>
<td>0.23</td>
<td>0.24</td>
</tr>
<tr>
<td>$^5\Lambda$He</td>
<td>2.11</td>
<td>$1.1 \times 10^{-2}$</td>
</tr>
<tr>
<td>$^6\Lambda$Li</td>
<td>0.08</td>
<td>$0.39 \times 10^{-3}$</td>
</tr>
</tbody>
</table>

9. – Systems with Strangeness $> -2$

Stable multi-strange hadronic matter is predicted on general grounds up to strangeness violating weak decay, and hyperons must contribute to neutron star matter, as the Pauli principle essentially assures their presence at some baryon density. Indeed it is expected
that roughly equal compositions of u, d, and s quarks, leading to a strangeness fraction \( f_s = -S/\Lambda \approx 1 \) and charge fraction \( f_Q = Z/\Lambda \approx 0 \) occurs in hadronic systems at densities between 2-3 times nuclear matter[107]. While the actual densities and binding energies are dependent on the unknown hyperon potentials, the stability of these calculations is robust.

A hadronic many body system can include a number of hyperons. Calculations using present baryon-baryon potentials indicate that such systems are particle stable, decaying weakly with a mean life approximately equal to that of the free \( \Lambda \). However, the decay is expected to proceed via the non-mesonic process, \( \Lambda N \rightarrow NN \). It has also been proposed that such systems could be more deeply bound, and better described by QCD rather than in terms of more conventional meson exchange potentials. These quark/glue systems are called “strangelets”, and may be even stable to both strong and weak decays [108]. The strangelet binding energy formula has been expressed by the form;

\[
B/A = -a_v + b_v y + a_s/A^{1/3} + a_c/A^{4/3} + a_x x^2 + a_y y^2
\]

In this expression, \( a_v \) is the volume energy, \( a_s \) is the surface energy, and \( a_c \) is the coulomb energy. The particle excess fractions are defined by;

\[
\begin{align*}
x &= (N - Z)/A \\
y &= (Z + N - \Lambda)/2
\end{align*}
\]

Here \( \Lambda \) is the number of \( \Lambda \)s.

Strangelets may only be produced in relativistic heavy ion collisions, and substantial numbers of strange particles are observed. It is possible that via coalescence, hot quark matter could cool into a strangelet[112]. However, no evidence for strangelets has yet been found, and there are no reliable estimates for their formation probabilities. Strangelet stability is predicted to lie near the line given by \( S + A = 0 \), or one strange quark per 3 total system quarks, with system density 1.5-2.0 times the nuclear density. It has been argued using the quark shell model, that systems with \( A \leq 6 \) should not be bound. This is consistent with the data, of course, but hardly definitive.

10. – Strange Nuclear Systems Formed in \( \bar{p} \) Reactions

Production of hypernuclei using anti-protons has been observed, and there are more detailed discussion of future possibilities [114] particularly with respect to double lambda systems in another session.
11. – Production of Nuclei having Baryons with Heavy Flavors

The possibility of stable charmed nuclei was raised[109] in 1977 after the discovery of charmed mesons. The \( \Lambda_c \) (2285 MeV) hadron is much more massive than the \( \Lambda_s \) (1116 MeV) and carries a unit positive charge. Its valence quark structure is \( \bar{c}q \), and it could be produced by meson-nucleon interactions through exchange or associated production of charm as is the strange \( \Lambda \). Theoretical studies have mainly focused on the stability of charmed nuclear matter. The issue involves the strength of the \( \Lambda_c \)-Nucleus potential, which must also include the Coulomb interaction. While a free \( \Lambda_c \) is stable with respect to the strong interaction, a \( \Lambda_c \) within a nucleus also undergoes the weak interaction \( \Lambda_c N \rightarrow NN \). A study of the s-shell charmed hypernuclei suggests that the \( ^3H \) is unbound, but \( ^2H \) should be bound [111].

Experimental investigation of charmed hypernuclei will be difficult. If mesonic production is considered, the short lifetimes of the D mesons, (\( \approx 1 \times 10^{-12} \)), make it impossible to use them as beams, and kinematics because of their large mass (1869 MeV), will require that they recoil with substantial laboratory momentum. However in a \( p(D, \pi)\Lambda_c \) reaction, the momentum transfer to the \( \Lambda_c \) is approximately a constant 450 MeV/s. The lifetime of the \( \Lambda_c \) is short, \( 2 \times 10^{-13} \), although this is sufficient long compared to the strong if not the electromagnetic interaction times, stable hadronic matter can form. Detection of charm hypernuclei must have fine granularity and high precision, and emulsion is ideally suited for this purpose. An hybrid emulsion detector has been previously used to study charm production [115], and could be easily adapted for hypernuclear work. A theoretical prediction of production cross sections would be useful. Experimental observations of charmed hypernuclei have been reported [110].

* * *

Any review article such as this is dependent in some measure on the work of others. I have attempted to recognize this work where possible. I also acknowledge the help of colleagues over the years in explaining to me various concepts of nuclear physics. It is to their credit where this paper has succeeded in producing a coherent presentation of this physics, and to my limitations where it has not. Hypernuclear physics has proved an exciting field for two generations, and I look forward to the next.

REFERENCES

[83] F. Doerrmann, International Conference on Electromagnetic Production of Strange
[86] K. Tanida and Z. Rudy, COSY Summer School, 1 (2003) 147