K Meson-Nucleus Interactions: Strangeness and Nuclear Matter

S. Kahana
Physics Department
Brookhaven National Laboratory
Upton, New York 11973

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S. Kahana

Physics Department
Brookhaven National Laboratory, Upton, New York 11973
U.S.A.

ABSTRACT

A brief review is provided of some straightforward K-nuclear and A-hypernuclear systems. A discussion of less straightforward speculations on H-dibaryons and strange quark matter by many authors, is also given.

The nominal subject of this talk is K-meson nucleus interactions. I will presume on the flexibility and tolerance of the organizers sufficiently to permit myself a selection of topics within the broader scope of the physics of a nucleus plus one or two units of strangeness. Thus I might touch on K*-nucleus scattering, A- or Σ-hypernuclei, singly and doubly strange dibaryons. I will in fact emphasize the latter rather exotic objects, going so far as to place them, after others, notably Witten, in the dark matter of the universe or in ultra-high energy objects arriving from the neighborhood of Cygnus X-3. I will, however, also pay some lip service to the few areas of K-meson physics with which I have had a more serious involvement.

To begin we note K+, K~ differ sharply in their interactions with nuclei, K~ meson may be absorbed through reactions like

\[ K^- + n \rightarrow \Lambda + \pi^- + \pi^- \]
\[ K^- + p \rightarrow \Lambda + \pi^0 + \pi^- \]
\[ + \Sigma^+ + 
\]
\[ \Sigma^0 \]
\[ \Sigma^0 \]
\[ \Sigma^0 \]
\[ \Sigma^0 \]
\[ \Sigma^0 \]
\[ \Sigma^0 \]

while K+ mesons cannot and are understandably more weakly interacting at low energies. Measured elastic scattering of K-mesons from nuclei are displayed in Fig. 1 in comparison to theory. Good fits to K-scattering are obtained with optical potentials, rather shallower for K+ than for K~ in both real and imaginary parts. Some such potentials are parametrized in Table 1. Of course, extensive information has been gathered experimentally on the elementary K+-nucleon amplitudes, and these amplitudes may be used to construct the K-nucleus interaction from, so to speak, first principles. A meson exchange model is permissible for K+N scattering. The model is dominated by ρ, ω (vector) or ϵ (scalar) mesons, the ω and η cannot contribute singly. For (K~,N), s-channel resonances play an important role and a multi-channel analysis is more appropriate. In
either case the effective K-nucleon interaction is short range and one can construct an optical potential from the approximation

\[ U(r) = V(r) + iW(r) = t\rho(r) \]  

where \( t \) is some appropriately averaged scattering matrix and \( \rho(r) \) is the nucleon density in a nucleus. Figure 2, shows a comparison of \( K^\pi \)-optical potentials obtained from (2) and an experimental knowledge of the basic two-body interaction. The calculation is for \(^{12}\text{C}\) scattering and clearly the \( K^+ \) potentials are shallower.

The processes in Eq. (1) are responsible for the production of hypernuclei. Extensive experimental\(^{5,6}\) and theoretical\(^{2}\) studies of the light mass, mainly p-shell, \( \Lambda \)-hypernuclei have uncovered considerable detail about the \( \Lambda-N \) interaction inside nuclei, some of it surprising. For example, it was inferred from experimental studies\(^{5}\) of \(^{16}\text{O}\) that the \( \Lambda \)-nucleus spin-orbit interaction was exceptionally small. A sophisticated shell-model analysis\(^{2}\) of the more appropriate odd-baryon nucleus \(^{13}\text{C}\) data allowed one to confirm this inference and in addition to pin down the two-body \( \Lambda-N \) interaction with some confidence. Some have argued\(^{7}\) that the small \( \Lambda-N \) spin-orbit force follows from the quark structure of the \( \Lambda \), i.e. the spin-orbit force is suppressed by the coupling of the up-down quarks to zero angular momentum and by the relatively high strange quark mass. However, meson-exchange models\(^{8}\) also lead to a small two-body spin-orbit.

Accurate treatments of the reaction mechanisms\(^{2}\) (shown in Fig. 3)

\[ K^- + ^{13}\text{C} \rightarrow ^{13}\text{C} + \Lambda + \pi^- \]  

using the elastic data for \( K^- \), \( \pi^- \) mesons near 800 MeV/c in momentum, are instrumental in producing the confidence in the extracted

---

**TABLE 1**

<table>
<thead>
<tr>
<th>Potential</th>
<th>Reaction</th>
<th>( V_0 ) (MeV)</th>
<th>( W_0 ) (MeV)</th>
<th>( r_0 ) (fm)</th>
<th>( a_0 ) (fm)</th>
<th>( \chi^2/N )</th>
<th>rms radius (fm)</th>
</tr>
</thead>
<tbody>
<tr>
<td>K1</td>
<td>( K^-+^{12}\text{C} )</td>
<td>24.4</td>
<td>41.4</td>
<td>1.075</td>
<td>0.375</td>
<td>0.31</td>
<td>2.36</td>
</tr>
<tr>
<td>K2</td>
<td>( K^-+^{12}\text{C} )</td>
<td>36.17</td>
<td>37.98</td>
<td>1.0</td>
<td>0.5</td>
<td>0.90</td>
<td>2.57</td>
</tr>
<tr>
<td>K3</td>
<td>( \pi^-+^{12}\text{C} )</td>
<td>0.9</td>
<td>50.9</td>
<td>0.926</td>
<td>0.44</td>
<td>2.5</td>
<td>2.32</td>
</tr>
<tr>
<td>K4</td>
<td>( K^-+^{40}\text{Ca} )</td>
<td>27.04</td>
<td>32.14</td>
<td>1.107</td>
<td>0.57</td>
<td>0.76</td>
<td>3.62</td>
</tr>
<tr>
<td>K5</td>
<td>( K^-+^{40}\text{Ca} )</td>
<td>26.34</td>
<td>28.06</td>
<td>1.134</td>
<td>0.55</td>
<td>0.76</td>
<td>3.63</td>
</tr>
<tr>
<td></td>
<td></td>
<td>23.57</td>
<td>18.69</td>
<td>1.182</td>
<td>0.49</td>
<td>0.71</td>
<td>3.63</td>
</tr>
</tbody>
</table>

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using the elastic data for \( K^- \), \( \pi^- \) mesons near 800 MeV/c in momentum, are instrumental in producing the confidence in the extracted
structure and force information (see Fig. 4). The elegant shell model analyses of D. J. Hillener at Brookhaven have contributed immeasurably to this entire subject recently. For example, these theoretical techniques as well as experimental refinements on \((K^-,\gamma)\) reactions throughout the \(p\)-shell, have broadened and deepened our knowledge of hypernuclear structure. \(\Sigma\)-hypernuclei, exhibiting what seem abnormally narrow states, have been of recent interest.

Measurements of \(A\)-hypernuclear lifetimes may present one with a unique means of exploring the nature of the four-fermion weak interaction

\[ A + N + N + N \]  

Suppression by the exclusion principle of the normal decay \(A + p\pi^+\) inside nuclei, projects the two-body mechanism, Eq. 4, into a dominant role, i.e. waiting the \(10^{-10}\) secs or so required for decay in effect turns off the strong interaction. The long range nature of \(4\) is not completely understood. Early results indicate a lifetime for \(A^0\) near that in free space.

Aside from these references to earlier work, I prefer to dwell on potentially more fundamental aspects of strange nuclear matter. As well as using single strangeness "changing" reactions, i.e. which implant one unit \(S=1\) on nucleons or nuclei, one can also exploit \((K^-,K^+)\) processes to introduce two units. Further than this one cannot go except perhaps in relativistic heavy ion collisions where many baryonic units of strangeness, albeit with small cross-section, can be created simultaneously. Indeed, one measurement of \(A\)-hypernuclear lifetime, \(\tau_A\), was made at Berkeley using a relativistic heavy ion beam. Rare examples of double \(A\)-hypernuclei exist in the literature. In particular an observation of a candidate reaction

\[ \Xi^- + ^{12}\text{C} + ^6\text{He} + X \ldots \]  

was made by Prowse. Both ensuing, sequential weak decays were visualized and analyses yielded for the \(\AA\) pairing energy, in the triply magic \(^6\text{He}\), a value close to \(4\) MeV. Such a value is consistent with our knowledge of the residual \(\AA\) forces likely to be acting between two \(A\)'s in nuclear \(S\)-states.

Despite this R. L. Jaffe has postulated the existence in nature of the intriguing \(H\)-dibaryon constructed from a flavour-symmetric quark configuration, \(u'd's^2\) in a single bag. Jaffe's reasoning rests on very general grounds: if the confining forces in QCD are flavour-blind, then the 1-gluon exchange symmetry energy, arising as in 2-flavour physics from the exclusion principle, will favour the \(\text{u'd's}^2\) six quark bag. Indeed, except for the not understood chiral breaking mass of the strange quark, this object would be absolutely stable relative to two nucleons. Of course, the strange mass puts the \(H\) somewhat higher in energy.
One can briefly sketch Jaffe's argument\(^1\). The 1-gluon exchange between s-wave quarks within the bag cavity is colour-magnetic in nature and gives an energy shift

\[
\Delta E_M = -\frac{a_c}{R} \sum_{i,j} \hat{\sigma}_i \cdot \hat{\sigma}_j \lambda_i^* \lambda_j^* h(m_i R, m_j R) \tag{6}
\]

where \(\hat{\sigma}_i, \lambda_i\) are the \(i\)th quark spin and colour matrices, \(a_c\) is the strong coupling constant at an appropriate scale, \(R\) is the bag size and \(h\) reduces essentially to products of intrinsic quark magnetic moments. Mass differences \(m_i\) amongst the quarks induce slight SU(3) flavour breaking in \(h\) and hence in (6), but to a good approximation.

\[
\Delta E_M = -\frac{a_c}{R} h(0,0) \sum_{i,j} \hat{\sigma}_i \cdot \hat{\sigma}_j \lambda_i^* \lambda_j^* \tag{7}
\]

In this form the energy may be evaluated using Casimir operators for the SU(6) group in which both colour SU(3) and spin SU(2) are embedded. The result is for a single \(N\)-quark bag coupled to a colour singlet,

\[
\Delta E_M = \frac{\alpha h}{R} \left[ 8N - \frac{1}{2} C_6 + \frac{4}{3} S(S+1) \right] , \tag{8}
\]

with for example \(C_6=90, S=1/2\) for the nucleon while \(C_6=42, S=3/2\) for the delta. Because of the non-abelian nature of the colour gauge theory the \(\Delta\) is heavier than the nucleon and the mass splitting is

\[
\delta E = 16 \frac{c}{R} \approx 315 \text{ MeV} \tag{9}
\]

which fixes the proportionality constant in (6). A more or less unambiguous prediction follows for the \(\Lambda\)-dibaryon. For six quarks in a single bag the double nucleon comes out as a highly excited resonance, perhaps evidence for the hard core. The large symmetry energy in Eq. (7) strongly favours the flavour singlet which possesses flavour-hypercharge, \(Y=S+\bar{S}=0\), and hence strangeness \(S=-\bar{B}=2\). For this object the colour-spin Casimir is \(C_6=144\) and (8), (9), (7) then yield a binding energy somewhat under 100 MeV. Jaffe\(^1\), including other SU(3) breaking finds \(M(\Lambda)\approx 2150\) MeV. More recent estimates\(^1\) reduce this value somewhat.

Such a deep binding seems to contradict the, admittedly sparse, experimental evidence on \(^6\)He. One might have expected the two \(\Lambda\)’s in s-states in \(^6\)He to fall into the well provided by the single \(H\) bag, during the relatively long \(\Lambda\) lifetime. Further higher energy, i.e. several GeV, experiments on

\[
p + p + K^+ + [X] \ldots \tag{10}
\]

at the AGS\(^1\), which could presumably have led to penetration of any hard core \(\Lambda-\Lambda\) barrier, have failed to uncover the \(H\) at the level of tens of nanobarns.
It is instructive to examine Jaffe's conclusions in a somewhat different context, within the framework of currently fashionable soliton models\textsuperscript{17,18} which respect SU(2) flavour symmetry, and potentially SU(3), as well. The solitons arise from a Yukawa coupling of quarks to the chiral scalar fields (SU(2))

\[
g\bar{\psi}(\dot{\sigma}(r) + i\dot{r} \cdot \pi(r)) \quad (11)
\]

For sufficiently strong coupling \(g\), one singly degenerate state in the \(B=1\) spectrum detaches itself and drops to negative energy while for the \(B=2\) spectrum two such states do this. The spontaneously symmetry-breaking of the solution

\[
\pi(r) = \hat{r}_f(r) \quad (12)
\]

governs this behavior.

Placing \(N_c=3\) quarks (SU(3) of colour) into the negative energy orbit is necessary to fill it and yields \(B=3(1/3)=1\) for the lowest state, a hadron. The hedgehog nature of the soliton form (12) was dictated by the \(\pi^*\tau\), SU(3) chirally invariant form of the interaction (11). Moreover, the degeneracy is determined by the conserved diagonal SU(2) generator

\[
\mathcal{G} = \mathcal{J} + \mathcal{t} \quad (13)
\]

which replaces the broken spin \(\mathcal{J}\) and \(\mathcal{t}\)-spin \(\mathcal{t}\). The lowest state has \(\mathcal{G}=0\) for which \(\mathcal{J}=\mathcal{t}=1/2\) in a single quark state, but which yields the nucleon and delta for 3 quarks.

If one extends the flavour group to include strangeness, the resulting SU(3) chiral structure suggests an interaction

\[
\sum_{i,a} \bar{\psi}_i Q^a(r)\lambda^a_{ij} \psi_j \quad (14)
\]

between the fundamental, three (SU(3)\(_f\)), quark fields \(\psi_i\) and an octet of pseudoscalar mesons \(Q^a(r)\). Three of the latter \(Q^{a=b=1,2,3}\) are still the \(\pi\)-meson but the other five presumably correspond to the \(K\) and \(\eta\) mesons. The most straightforward generalization of (12) is then

\[
Q^i(r) = \hat{r}^i \quad (15)
\]

\[
Q^a(r) = \hat{r}^a \hat{r}^j \frac{1}{3} \delta_{ij}
\]

i.e. an SU(3)\(_f\) hedgehog. Now, however, to couple angular momentum to the 3-d representation of SU(3) flavour we would require an \(sl/2\), \(d3/2\) (or \(pl/2, p3/2\)) pair of quark orbits.
Simple counting

\[ 2(1/2)+1 + 2[3/2]+1 = 6 = 2 \cdot 3 \]

indicates this orbit is doubly degenerate which in turn implies the lowest SU(3) flavour-symmetric object has \( B=2 \) rather than \( B=1 \) and at the same time \( S=-2 \). This is again the \( H \)-dibaryon. The SU(3)\(_f\) symmetry is, of course, broken by the current mass of the strange quark, the \( B=2 \), \( H \) is no longer absolutely stable but is still favoured relative to other 6-quark solitons.

Further Experimental Search for the \( H \)

We have indicated some experimental evidence\(^{13, 16}\) exists, apparently, contradicting the predictions for the \( H \). Aerts and Dover\(^{19}\) have suggested alternative reactions for dropping two units of strangeness into an elementary system. One of these is shown schematically in Fig. 6. Results of cross-section calculations for this reaction

\[ K^- + \text{^3He} \rightarrow \text{H} + \text{n} \quad (16) \]

are shown in Fig. 5. The pp mechanism \((10)\), performed at a few GeV incoming energy, can overcome an anticipated hard core in the baryon-baryon short range force. This barrier might have explained a failure to achieve the \( H \) in \( ^6\text{He} \). Aerts and Dover evaluate the required quark fusion in terms of a form factor involving the proton and \( \Xi^- \) relative momentum in Fig. 6.

\[
\Gamma(k_p, k_\Xi) = \exp\left[-\frac{R_H^2}{12} (k_p - k_\Xi)^2\right] \cdot f(R_p, R_H, R_\Xi) \quad (17)
\]

This form factor highly favours \( H \) formation in the \( K^- + \text{^3He} \) reaction where \( k_p - k_\Xi \) is small. In contrast the \( AA \) quark fusion occurring in reaction \((10)\) is at a large relative momenta. However, the enhancement found by Aerts and Dover may be reduced by the hard core they have neglected.

An experiment using improved \( K^- \) beam at the AGS and exploiting \((16)\) or the possibility of \( H \) formation from \( \Xi^- \)-atoms, has been proposed\(^{20}\). Estimates using Ref. \((19)\) indicate a high production of intermediate cascades \( \Xi^- \) and hence a good likelihood for success in the search. If the present negative evidence for the \( H \) should persist, the simple assumptions behind Jaffe's analysis will be in question. Serious implications for QCD will follow. One way out is to allow the confining force to be flavour dependent. No evidence for such behaviour is evident, for example, in heavy quarkonium.

Stable, Strange Quark Matter

Considerable attention has attached to the subject of quark matter\(^{21}\) consisting of gluons and \( u, d \) quarks. This matter is expected, for vanishing chemical potentials, to appear above a
critical temperature $T = T_c$ after a deconfining phase transition from ordinary hadron matter. Since this temperature is anticipated near $T_c \approx 200$ MeV strange quarks will also play a strong role. Corresponding phase transitions are expected for baryon densities ten times normal saturation density (i.e. a few $x 10^{15}$ gm/cc), at low temperature. The quark matter systems will not exist at zero pressure, requiring a high pressure environment to stabilize. With a sufficiently high fermi momentum for these systems it will become favorable to include the third species of quark. The fermi momentum is thereby reduced and the mass difference $m_s - m_u, d$ compensated. This symmetry effect, arising from fermi-statistics, is familiar from nuclear physics. Indeed Freedman and McLerran\textsuperscript{22} had postulated absolute stability for a sufficiently large chunk of strangeness rich quark matter. Kerman and Chin\textsuperscript{23} have discussed a similar possibility.

Witten\textsuperscript{1} has pushed this concept further and has speculated on the nature of such matter its role in cosmology and in astrophysics. Strange quark matter, which is completely SU(3) symmetric is electrically neutral; the H is a microscopic beast of this type. Freedman and McLerran, however, find for one model a ratio of up to down to strange abundances of 2:3:1. These ratios generate chemical potentials $\mu_u / \mu_d = (2/3)^{1/3}$ and $\mu_q - \mu_u = 50$ MeV. Their system is notelectrically neutral and will repel normal charged nuclei but not neutrons. It then is likely neutron stars are in danger of burning reasonably quickly into such matter, provided some of this matter is produced at the center of the remnant left after a supernova and provided the temperatures throughout the neutron star are sufficiently high.

I end this imaginative sojourn by the authors I have cited, with even wilder speculations of other authors. Very recently, first the proton-lifetime SUDAN group\textsuperscript{24} and then the NUSEX collaboration\textsuperscript{25} turned their attention to high energy $\mu$-mesons appearing in their detectors. It is known\textsuperscript{26} that the binary neutron star Cygnus X-3 is a hot source of X-rays and of much more energetic $\gamma$-rays. The SUDAN group have correlated the $\mu$'s they detect with the 4.8 hour period, and even the phase, of Cygnus X-3. Detailed analysis yields several facts:

(1) The flux of $\mu$'s, with earth surface energies of several hundred GeV, cannot be produced by primary $\gamma$'s.

(2) The $\mu$'s are not neutrinos.

(3) The primaries must be neutral (spatial dispersion).

(4) The primaries probably have a mass not much larger than 1-2 GeV (time dispersion).
(5) The $\mu$'s are correlated with a near extinction of the neutron star by its companion.

These points have led some to invoke Witten's scenario and to speculate quark matter is arriving here from the Cygnus X-3 neutron star after collision with the binary companion. Others invoke both Witten and Jaffe's arguments to conclude quark matter from the neutron star produces $H$'s in collisions with the companion. The $H$'s then arrive at the Earth, the only speculation apparently consistent with fact (4). In this latter scenario, however, the $H$ is given a binding energy in excess of 180 MeV and is consequently stable against single weak decay.

References


11. Recent measurements at the AGS at BNL by a Carnegie-Mellon University group have extracted lifetimes for Mass 12 and 13 $A$-hypernuclei of near 200 p-secs, somewhat faster than the free lifetime $t \simeq 260$ p-sec.


20. A search for the H-dibaryon using the reaction in Ref. 19 has been proposed at the AGS by the CMU group. This experiment would be part of a broader study using an unproved K beam at the AGS. (Proposal for a 1-2 GeV/c Beam Line for Hypernuclear and Kaon Research, R. Chrien et al.; L. McLerran and B. Svetitsky, Phys. Rev. B24 (1981) 450; J. D. Bjorken, Phys. Rev. D27 (1983) 140.


Fig. 1 K-Nucleus Elastic Scattering. An optical fit to the 800 MeV/c data for $K^-$ (and $\pi^-$) mesons incident on $^{12}$C. The data is from Marlow et al.,$^3$ while the curves labelled K1, K2, $\pi$1 refer to the Wood-Saxon potential parameters given in Table 1.
Fig. 2 Theoretical $K^\pm$ optical potentials. Real and imaginary parts $V(r)$ and $W(r)$ of the $K^\pm + ^{12}\text{C}$ optical potentials at 500 MeV/c in the $t\rho(r)$ approximation, as calculated in Ref. 4.
Fig. 3 The $^{13}$C theoretical spectrum. The $^{12}$C($0^+, 2^+$) $\times p_A$ spectrum for an interaction independent of $\sigma_A$. Parent core states in $^{12}$C, constructed configurations in $^{13}$C; their spatial symmetries and total orbital angular momentum are indicated above the figure. Substitutional states dominating the strong experimental states at 10 MeV and 16 MeV are indicated by an asterix.
Fig. 4 Comparison of measured angular distributions for $K^- + ^{13}C \rightarrow \pi^- + ^{13}C$ with absolute theoretical predictions. The optical potentials used in the relativistic DWBA analysis are those from Fig. 1 and Table 1. The data is very well described.
Fig. 5 Differential cross-sections for H-dibaryon production on a $^3$He target, from Ref. 19. Both the forward differential cross-section at 1.9 GeV as a function of the H mass, and the laboratory cross-section for a K$^+$ at 0° as a function of the K$^+$ momentum $k'$, are shown.

Fig. 6 Schematic Feynman diagrams for the production of a H-dibaryon in the K$^+$$^3$He reaction. Also shown is one contribution to the background. This figure is from Ref. 19. If the H is bound, the background and signal separate.