STRANGE QUARKS IN LOW ENERGY HADRONIC INTERACTIONS

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The dynamics of the K^+N-interaction is calculated in the constituent quark model with QCD-residual forces. The obtained S-wave interaction agrees rather well with the experimental data. We also discuss the status of spin-orbit forces in KN, Λ-nucleus and Σ-nucleus interactions.

I. Introduction

Strong interactions are independent of flavor. The forces mediating the strong interaction are colored vector-gluons which carry no flavor. The masses of elementary particles and their interactions nevertheless show strong flavor-dependence: e.g., the K-meson with $m_K \approx 500$ MeV is much heavier than the π-meson with $m_\pi = 137$ MeV. The Lagrangian density of QCD includes this flavor asymmetry by a nontrivial mass matrix M^{ij}, i.e.

$$L(x) = -1/4\, F^a_{\mu\nu}(x) F^{a\mu\nu}(x) + \bar{\psi}^{i,\alpha}(x)\,(i\gamma^u D^{\alpha\beta}_\mu \delta^{ij} - M^{ij})\psi^{j,\beta}$$

$$D_\mu^{\alpha\beta} = \partial\mu - ig \frac{\lambda_a^{\alpha\beta}}{2} A^a_\mu \; .$$

For only three flavors the mass matrix M^{ij} has the form:

$$M^{ij} = \begin{pmatrix} m_u & 0 & 0 \\ 0 & m_d & 0 \\ 0 & 0 & m_s \end{pmatrix} \quad \text{with } m_s \approx \frac{m_u+m_d}{2} + 0.2 \text{ GeV}$$

In this picture the chiral SU3 x SU3 symmetry is broken due to the large strange quark mass. We therefore believe that chiral dynamics plays a less important role in strange hadron-hadron interactions. If we loosely identify chiral dynamics with pion or pioncloud effects, we would expect that long-range meson exchanges are not distorting the short-range QCD-picture in strange quark interactions. Indeed the situation is very clear for K^+N-scattering which has clear advantages to NN-scattering. For the K^+N-interaction

(i) one-pion exchange is not allowed, because the coupling of three pseudo-scalars has to vanish;

(ii) two-pion exchange is suppressed, since the $KK^*\pi$ coupling is smaller than the $N\Delta\pi$ coupling.

The K^* has half the width of the Δ in spite of the larger phase space in $K^* \to K\pi$ compared to $\Delta \to N\pi$. In addition the t-channel $NN \to NN$ amplitude has the square of $(N\bar{N} \to \pi\pi)$ amplitudes as weight functions, whereas in $K^+N \to K^+N$ the $\bar{K}K \to \pi\pi$ and $\pi\pi \to N\bar{N}$ amplitudes may have different signs. Hence the quark model alone may give rather realistic phase shifts for KN-scattering. Another advantage of the quark description in the KN-system is the isospin selectivity. Different isospins show complementary features in the KN-cluster model. The $I = 0$ channel is determined by quark-gluon exchange only, whereas the $I = 1$ channel has contributions to the scattering from quark rearrangement and quark-gluon exchange. Therefore, in the KN-system we are able to investigate the relative importance of the QCD-residual interaction and of quark antisymmetrization.

It is sometimes incorrectly stated that the repulsive NN-interaction is a consequence of quark antisymmetry. This is not true. Quark rearrangement alone leads to a vanishing interaction for NN-scattering. Only antisymmetrization together with the QCD-residual force, i.e. quark-gluon exchange, gives a repulsive short-range NN-interaction. Historically this work on strange quarks has been triggered by the experimental observation of Λ-hypernuclear spectra. In the meantime also Σ-hypernuclei have been seen. I will discuss their preliminary interpretation in the context of quark models. Let me start, however, with the KN-interaction. The understanding of this system would represent an important step forward towards a theory of hypernuclei. For a detailed discussion I would like to refer to ref. [1].

II. K^+N-scattering

In the constituent quark model all the complications and subtleties of non-perturbative QCD are parametrized by the constituent quark mass and a confining potential. The Hamiltonian has the form

$$H = \sum_{i=1}^{4} \frac{p_i^2}{2m} + \frac{p_5^2}{2m_s} + \sum_{i<j} \mathcal{V}_{ij}^{conf} + \sum_{i<j} \mathcal{V}_{ij}^{QCD} \qquad (1)$$

where the confinement is included via the two-body interaction

$$\mathcal{V}_{ij}^{conf} = -\frac{D}{2} \lambda_i \lambda_j \; |x_i - x_j|^2 \tag{2}$$

and the QCD-residual interaction is the non-relativistic limit of the one-gluon exchange potential (without LS and tensor terms for the S-wave interaction).

$$\mathcal{V}_{ij}^{QCD} = \frac{\alpha_s}{4} \lambda_i \lambda_j \; \{\frac{1}{|x_i - x_j|} - \frac{2\pi}{3} \frac{1}{m_i m_j} \; \sigma_i \sigma_j \; \delta(x_i - x_j) - \frac{\pi}{2} (\frac{1}{m_i^2} + \frac{1}{m_j^2}) \; \delta(x_i - x_j)\} \tag{3}$$

We use the following parameters: the constituent quark mass is m = 0.3 GeV, the oscillator parameter D is adjusted to reproduce the p-wave nucleon excitation at 1.535 GeV. The QCD-coupling constant $\alpha_s = 0.7$ is obtained from the ΔN-splitting. The strange quark mass is related to the Λ, Σ and N mass difference.

The Hamiltonian of eq. (1) has as asymptotic states the kaon and the nucleon. Color confinement does not act asymptotically because of the color neutrality of the hadronic clusters. It is natural (but not unique) to choose a trial 5-quark wavefunction which takes into account the hadronic cluster states ϕ_N and ϕ_K.

$$\psi = A \int c(R) \; \psi_{\vec{R}} \; d^3R \tag{4}$$

with

$$\psi_{\vec{R}} = \phi_N(\xi_1 \xi_2) \; \phi_K(\xi_3) \; \chi_{\vec{R}}(r) \; \phi_{cm}(R_G) \; \chi_{KN}^{SIC} \tag{5}$$

and $A = 1 - P_{14} - P_{24} - P_{34}$.

The trial wavefunction of eqs. (4,5) is a superposition of hadron clusters at relative distances R. In the interaction region it is advantageous to choose $\chi_{\vec{R}}(r)$ as a Gaussian function peaked around this generator coordinate R and let the center of mass R_G free to oscillate. With this choice the wavefunction $\psi_{\vec{R}}$ can be reexpressed in terms of quark coordinates $x_1, \ldots, x_5$ (see table 1). ψ_R contains two shell model wavefunctions peaking around $-(1 + K)/(4 + K)\; \vec{R}$ and $+ 3\; \vec{R}/(4 + K)$, where $K = m_s/m$ is the mass asymmetry.

$$\psi_R(x_1, x_2, \ldots x_5) = A \{ \prod_{i=1}^{3} (\frac{2\Omega}{\pi})^{3/4} e^{-\Omega(\vec{x}_i + \frac{1+K}{4+K}\vec{R})^2} (\frac{2\Omega}{\pi})^{3/4} e^{-\Omega(\vec{x}_4 - 3\vec{R}/(4+K))^2} (\frac{2K\Omega}{\pi})^{3/4} e^{-\Omega K(\vec{x}_5 - \frac{3\vec{R}}{4+K})^2} \} \tag{6}$$

Note the heavier strange quark is more localized at the K-cluster because of $K \approx 1.7$ it will not participate as much in peripheral partial waves.

The key (ref. [1]) to the solution of the time-independent Schrödinger equation

$$(H - E)\ \psi = 0 \tag{7}$$

is the projection of this equation onto a variation of the many quark wavefunction (4) with respect to $c(\vec{R})$:

$$\int \delta\psi\ (H - E)\ \psi\ d^3r d\tau^{12} = 0 \tag{8}$$

$$d\tau^{12} = \prod_{i=1}^{3} d^3\xi_i\ d^3R_G$$

For large distances r the relative motion of the two clusters is described by a free wavefunction. Hence by applying Green's formula to the integration over d^3r we obtain

$$\begin{aligned} \int_G d^3r \quad d\tau\ [\delta\psi (H - E)\ \psi - \psi (H - E)\delta\psi\ ] &= \\ - \frac{1}{2\mu} \int_{\partial G} d^2s \quad d\tau\ (\delta\psi \frac{\partial\psi}{\partial n} - \psi \frac{\partial\delta\psi}{\partial n}) &= - \frac{1}{2\mu k}\ \delta S \end{aligned} \tag{9}$$

where $\mu = \frac{3(1+K)}{4+K}\, m$ is the reduced mass of the KN-system, G is a bounded region in R^3, ∂G its surface, and $\partial/\partial n$ the normal derivative. Using a superposition of incoming and outgoing spherical waves the r.h.s. of eq. (7) is proportional to the variation of the S-matrix δS. The Schrödinger equation can be reformulated as a variational equation $\delta I = 0$ with

$$I = \frac{1}{2\mu k}\ S - \int_G d^3r \quad d\tau\ \psi (H - E)\psi \tag{10}$$

which leads to a linear system of equations for a discretized set of coefficient $c(R_i) = c_i$ in eq. (4). The whole physics is contained in the expectation values of $(H - E)$ between cluster states located at distances $\underset{\sim}{R}$ and $\underset{\sim}{R}'$.

$$H(\underset{\sim}{R}',\underset{\sim}{R}) - N(\underset{\sim}{R}',\underset{\sim}{R})\ E = \int \psi_{\vec{R}'}\ (H - E)\ A\ \psi_{\vec{R}}\ \ d\tau \tag{11}$$

In general there are direct and exchange parts for each kernel from the two parts of the antisymmetrizer $A = 1 - 3\ P_{34}$. These are represented in fig. 1. A cross symbolizes a kinetic energy term $p_i^2/2m$, the dotted lines stand for a quark-quark interaction.

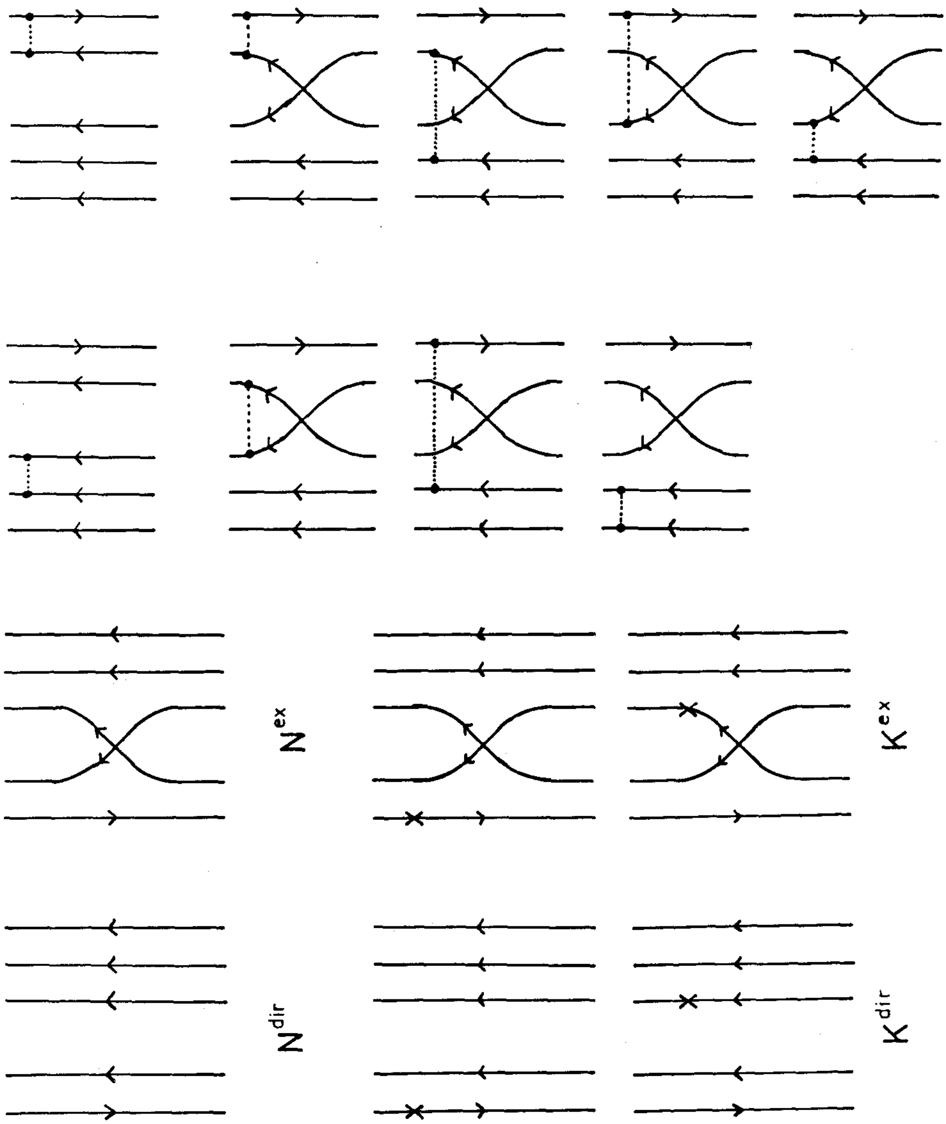


Fig. 1 Contributions to the normalization kernel N($\underset{\sim}{R}$,R'), kinetic energy kernels K($\underset{\sim}{R}$,$\underset{\sim}{R}$') and potential kernel V(R,R'), where H(R,R') = K(R,R') + V(R,R').

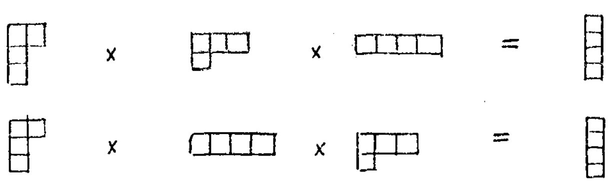
The spin color overlap factors are such that

$$< \chi^{SIC} \mid P_{34}^{SIC} \mid \chi^{SIC} > = 1/9\ \delta_{I1}\ , \tag{12}$$

i.e., the exchange operator vanishes for the KN-cluster function in the I = 0 state. Therefore quark exchange alone from N^{ex} and K^{ex} (fig. 1) does not contribute to the I = 0 state. The effect of quark exchange is very strong, however, in the I = 1 state, as one can see from a solution of the problem without QCD-residual interaction, i.e. $\alpha_s = 0$.

In fig. 2 we show the phase shift obtained from the calculation in the I = 1 channel with $\alpha_s = 0$. The theoretical phase shift already reproduces the data (crosses).

The origin of this repulsion can be visualized by regarding the spin-isospin symmetry of the wavefunction. Because of the color symmetry [2, 1, 1] of the 4 light quarks only the SU(4) $[3, 1]_{IS}$ states are allowed for a spatially symmetric wavefunction [1 1 1 1]. The SU(4) state $[4]_{IS}$ has to be associated with a spatially excited 4-quark state:



(13)

Color- $\quad$ $SU(4)_{IS}$- $\quad$ spatial $\quad$ total wavefunction

In line with this symmetry argument the quark exchange due to antisymmetrization generates different potentials for the two SU(4) states: the sign of the potential is positive for the $[4]_{IS}$ and negative for the $[3, 1]_{IS}$ spin-isospin wavefunction. Denoting by P_{34}^{SI} the exchange operator of particle 3 and 4 acting on the spin-isospin wavefunction one obtains

$$\begin{aligned} &< [4]_{IS} \mid P_{34}^{SI} \mid [4]_{IS} > = 1 \\ &< [3,\ 1]_{IS} \mid P_{34}^{SI} \mid [3,\ 1]_{IS} > = -1/3 \end{aligned} \tag{14}$$

The I = 0 K-N wavefunctions can be expressed as a superposition of SU(4) supermultiplet states coupled to a strange antiquark

$$\begin{aligned} &\mid KN;\ I = 0,\ l = 0 > = -1/2([4]_{00}\ \bar{s}) + \frac{\sqrt{3}}{2}\ ([3,\ 1]_{01}\ \bar{s}) \\ &\mid KN;\ I = 1,\ l = 0 > = \frac{1}{\sqrt{2}}\ ([4]_{11}\ \bar{s}) - 1/2([3,\ 1]_{10}\ \bar{s}) - 1/2([3,\ 1]_{11}\ \bar{s})\ . \end{aligned} \tag{15}$$

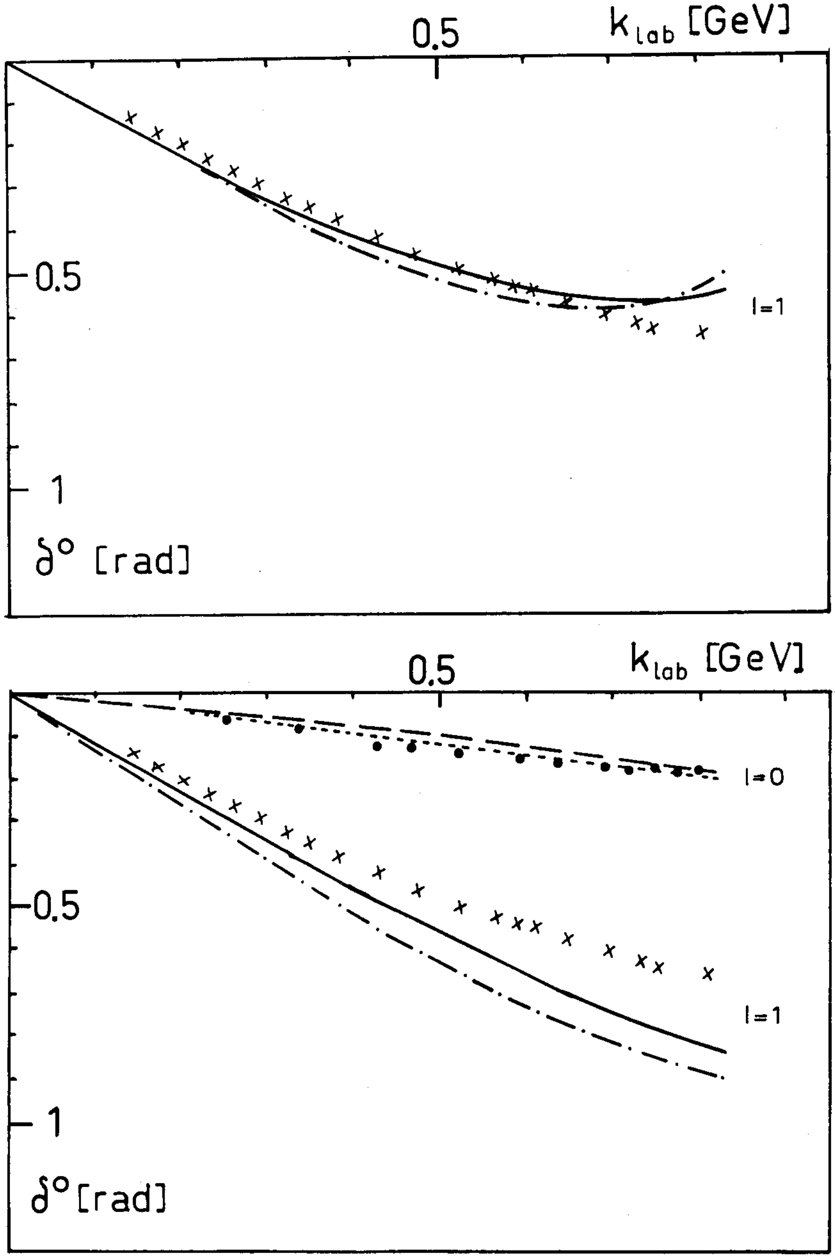


Fig. 2b Theoretical evaluation of the KN-S-wave phase shifts for I = 0 and I = 1 and experimental data from ref. [8]. The upper curves for each isospin correspond to the smaller cluster size $\langle r^2\rangle^{1/2} \approx 0.43$ fm compared to $\langle r^2\rangle^{1/2} \approx 0.47$ fm.

Combining eqs. (14) and (15) we see that for the $I = 0$ state the attractive $[3, 1]_{IS}$ and repulsive $[4]_{IS}$ contributions of the exchange potential cancel, whereas for the $I = 1$ state the repulsive $[4]_{IS}$ contribution dominates. This explains the strong repulsion in the $I = 1$ channel which is produced by the antisymmetrization of the 4 light quarks.

In ref. 1 we have calculated a local approximation to the Hill-Wheeler equation; it gives an r-dependent effective mass which acts like a repulsive shell prohibiting the two hadrons to approach. Including the QCD-residual interaction in addition (fig. 2b) we obtain weakly repulsive $I = 0$ phase shifts in agreement with the data, but somewhat too large phase shifts for the $I = 1$ state. We hope that a coupled channel calculation with the K^* and Δ-cluster states can improve these results. In general the agreement between theory and experiment is rather good. Comparing to NN-phase shifts one wonders what is the major difference between the two cases? We think that the attractive 2π-exchange which is sizable in NN-scattering but weak in KN-scattering is causing the different behavior. It remains a challenging task to incorporate such non-perturbative QCD effects in the NN-problem. On the other hand the KN-system has the good feature of weak long-range forces. Thereby it presents a much better candidate to study the properties of the perturbative QCD-interaction for multiquark states. We have also calculated the P-waves and compared to an average of the experimental amplitudes taken in such a way that a weak spin-orbit force cancels. There is no correlation between theory and experiment, however. It is therefore interesting to see whether agreement between theory and experiment can be obtained by including QCD-residual spin-orbit and tensor terms, or whether exotic $(qqqq\bar{q})$ configurations play an important role in the P-waves.

III. Spin-Orbit Interactions

Previous investigations have shown the P-wave KN-interaction can be fitted by a pure spin-orbit interaction obtained from quark-additivity in the limit of strong SU3-breaking.

$$\mathcal{V}_{LS}^{KN} = 30 \text{ MeV fm}^2 \frac{1}{r}\frac{\partial}{\partial r}\frac{\tilde{\rho}(r)}{\tilde{\rho}_o}\Big(\sum_{Q\in N} \vec{s}_Q(1 + \vec{\tau}_K\vec{\tau}_Q)\Big)(\Sigma\vec{l}_Q) \tag{16}$$

In the expression of eq. (16) we find the familiar Thomas-form of the spin-orbit potential with a strength adjusted to the baryon-nucleus spin-orbit interaction. The exchange of the quarks leads to the isospin exchange operator $1/2(1 + \tau_K\tau_Q)$, since the isospin of the K is equal to the isospin of the exchanged light quark. Taking the expectation value of eq. (16) with respect to the nucleon we find

$$\mathcal{V}^{KN}_{LS} = 30 \text{ MeV fm}^2 \frac{1}{r}\frac{\partial}{\partial r}\frac{\tilde{\rho}(r)}{\tilde{\rho}_o} \underset{\sim}{s}\cdot\underset{\sim}{l}(1 + 1/3\ \underset{\sim}{\tau}_K\underset{\sim}{\tau}_N) \qquad (17)$$

The 5/3 appearing in eq. (17) is the same as in the naive quark model prediction for the axial coupling constant g_A. The normalized quark distribution has been assumed to be $\Theta(R_o - r)$ with $R_o \simeq 0.5$ fm. This model fits the data quite well (ref. [2]). It is built in analogy to the baryon-nucleus spin-orbit interaction which is obtained in ref. [3] by adding quark-nucleon interactions in the nucleus and averaging over the quark-wavefunction of the valence baryon.

$$\mathcal{V}^{B}_{LS} = 90 \text{ MeV fm}^2 \frac{1}{r}\frac{\partial}{\partial r}\frac{\rho(r)}{\rho_o}\ 1/n \sum_{u,d}^{n} \underset{\sim}{s}_Q\cdot\underset{\sim}{l}_Q \qquad (18)$$

giving the spin-orbit potentials for the N, Δ, Λ and Σ-nucleus interaction of table 2.

Two questions have to be discussed:

(1) Are the recently measured Σ-hypernuclear data in agreement with this phenomenological model based on quark-additivity?

(2) Can quark-additivity be derived in the constituent quark model from the ΛN and ΣN spin-orbit interactions?

Concerning the first question we have come to the following conclusion. If the peak in the spectrum on $^{12}C(K^-,\pi^+)^{12}C_\Sigma$ coincides with the lower energy ($m_{Hy} - m_A$ smaller) peak in the $^{16}O(K^-,\pi^+)^{16}O_\Sigma$-spectrum, then a residual interaction of the form

$$\mathcal{V}^{\Sigma N} = C_o\ \delta(\underset{\sim}{r}_\Sigma - \underset{\sim}{r}_N)\ [f_{\Sigma N} + f'_{\Sigma N}\ \underset{\sim}{\tau}\underset{\sim}{T} + g_{\Sigma N}\ \underset{\sim}{\sigma}_\Sigma\underset{\sim}{\sigma}_N + g'_{\Sigma N}\ \underset{\sim}{\sigma}_\Sigma\underset{\sim}{\sigma}_N\ \underset{\sim}{\tau}\underset{\sim}{T}] \qquad (19)$$

has to give very small mixing of $|3/2\ 3/2^{-1}>$ and $|1/2\ 1/2^->$ states, i.e.

$$< \Sigma_{3/2}\ p^{-1}_{3/2}\ \ |\mathcal{V}^{\Sigma N}|\ \Sigma_{1/2}\ \ p^{-1}_{1/2} > \quad \approx 0 \qquad (20)$$

which puts strong constraints on the coefficients f,f',g,g'. More importantly the peak corresponding to the lower mass in O^{16} has to be the $\Sigma_{3/2}\ \ p^{-1}_{3/2}$ state. The necessary spin-orbit strength of the Σ then is $\approx$ 5/3 as big as the nucleon spin-orbit strength. The Coulomb energy differences do not change the above statement. The argument is very simple, since we know the undisturbed mass $m^o_{3/2}$ in C^{12} and the sum of the masses $(m_{3/2} + m_{1/2})$ in O^{16} we can deduce the undisturbed level $m^o_{1/2}$

$$m^o_{1/2} = m(\text{lower}) + m(\text{upper}) - m^o_{3/2} \approx m(\text{upper}) \qquad (21)$$

No residual interaction can invert the assignment m(lower) = $m_{3/2}$ and m(upper) = $m_{1/2}$, having taken into account the information from C^{12}. For the residual interaction above

M. Sommermann [4] gets the relation of the coefficients

$$f_{\Sigma} - f'_{\Sigma} + 3g_{\Sigma} - 3g'_{\Sigma} = 0 \qquad (22)$$

which is compatible with an extra-polation of the NN-Migdal-interaction using quark-additivity

$$f_{\Sigma N} = 2/3\ f_{NN}\ ,\ f'_{\Sigma N} = 2\ f_{NN}\ ,\ g_{\Sigma N} = 4/3\ g_{NN}\ ,\ g'_{\Sigma N} = 4/5\ g'_{NN} \qquad (23)$$

and modifying the I = 1/2 channel because of the presence of the $\Sigma N \rightarrow \Lambda N$ channel by a renormalization factor Z. The new Migdal constants $\tilde{f},\tilde{f}'$ are given by

$$\tilde{f}_{\Sigma N} = 2/3(f_{\Sigma N} + f'_{\Sigma N}) + 1/3\ Z(f_{\Sigma N} - 2f'_{\Sigma N})$$
$$\tilde{f}'_{\Sigma N} = 1/3(f_{\Sigma N} + 2f'_{\Sigma N}) + 1/3\ Z(-f_{\Sigma N} + f'_{\Sigma N}) \qquad (24)$$

and $\tilde{g},\tilde{g}'$ obey corresponding equations. The Σ-well depth of half the well depth of normal nuclei corresponds to $Z \approx 1.5$.
It must be admitted, however, that the isospin splitting of "$\Sigma^{+}p^{-1}$ and $\Sigma^{0}n^{-1}$" states observed in the $^{12}C(K^{-},\pi^{-})$ reaction would come out too big with the above coefficients.

Taken the above interpretation with two peaks in the ^{16}O-spectrum there is no way other than to conclude that the Σ-hypernuclei have a large spin-orbit splitting. The meson exchange parametrizations have to take into account stronger SU3-breaking and/or include the damping of two-pion exchange also in the real part, since it does not lead to the correct width of the Σ-hypernuclear states.

The second question about a microscopic derivation of the additive quark model for the spin-orbit interaction in the constituent quark model is rather interesting, too. Recently two papers [5, 6] have appeared which have calculated the spin-orbit BN-forces in the quark model. In ref. [5] the spin-symmetric part of the spin-orbit one-gluon exchange is used only. It gives a rather good spin-orbit NN-interaction compared to phenomenological interactions. The most general spin-orbit baryon-nucleon potential has two contributions: a symmetric $L(\underset{\sim}{s}_B + \underset{\sim}{s}_N)$ and antisymmetric $L(\underset{\sim}{s}_B - \underset{\sim}{s}_N)$ part. For the ΛN-interaction both terms are big but their sum almost cancels. Also for the ΣN-potential the antisymmetric part is considerable. The final ratios are given in table 2.

The work of ref. [6] treats both the symmetric and antisymmetric spin-orbit interactions from one-gluon exchange. In addition it includes possible confinement terms which help to solve spin-orbit splittings of mesons but actually have a very bad effect on baryon spin-orbit splittings. The calculations are done in the flavor SU3-limit and yield an almost as strong ΛN and ΣN-potential. Both references correctly criticize the relativistic calculation of ref. [3] which only takes into account one possible diagram. However, only the outcome of a calculation with SU3-breaking will answer the question of whether the nuclear spin-orbit splitting is a remnant of the quark-quark interaction or not. We are working on this question for the P-wave KN-interaction which is so similar to the baryon-nucleus case. Hopefully then, one can also settle the question of the baryon-nucleus spin-orbit potential.

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Tables:

table 1 $K = m_s/m$

quark coordinates	masses	cluster coordinates
$\underset{\sim}{x}_1$	m	$\underset{\sim}{\xi}_1 = \underset{\sim}{x}_1 - \underset{\sim}{x}_2$
$\underset{\sim}{x}_2$	m	$\underset{\sim}{\xi}_2 = x_3 - \frac{x_1+x_2}{2}$
$\underset{\sim}{x}_3$	m	$\underset{\sim}{\xi}_3 = x_4 - x_5$
$\underset{\sim}{x}_4$	m	$\underset{\sim}{r} = \frac{1}{3}(x_1+x_2+x_3) - \frac{1}{1+K}(x_4+Kx_5)$
$\underset{\sim}{x}_5$	Km	$\underset{\sim}{R}_G = \frac{1}{4+K}(x_1+x_2+x_3+x_4+Kx_5)$

table 2

Spin-orbit potential for the N,Δ,Λ and Σ nucleus and the KN(I = 1) and KN(I = 0) systems as $V_{LS} = \omega \times 30$ MeV fm^2 $(1/r)(\partial/\partial r)[\rho(r)/\rho_o]\underset{\sim}{S}\cdot\underset{\sim}{L}$.

System						
Spin-orbit: potential (ω)	N	Δ	Λ	Σ	KN (I=1)	KN (I=0)
Additive quark model	1	1	0	4/3	8/3	-4
Constituent quark model (ref.5)	1	?	0.21	0.55	?	?