Roles of Quark Degrees of Freedom in Hypernuclei

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The quark model description of the hyperon nucleon forces, especially the antisymmetric spin-orbit forces, is studied from the spin-flavor SU(6) and the flavor SU(3) symmetry point of view. It is pointed out that the quark exchange interaction predicts strong antisymmetric spin-orbit force between the hyperon and nucleon.

1. Introduction

In this talk, I cover the following four subjects.

1. Quark model description of the baryon-baryon interaction including hyperons\textsuperscript{[1]}.  
2. Antisymmetric spin-orbit force from the SU(3) flavor symmetry point of view\textsuperscript{[2]}.
3. Description of weak $\Lambda N \rightarrow NN$ interaction in the direct quark mechanism\textsuperscript{[3]}.
4. Magnetic moments of light hypernuclei and contributions of $\pi$, $K$ exchange currents\textsuperscript{[4]}.

In this report, I discuss the first two subjects in detail and leave the others to references given above.

2. Baryon-Baryon Interactions

Recent experimental activities in hypernuclear physics provide us with high quality data of production, spectra and decays of hypernuclei. The accumulation of such data accelerates quantitative analyses of the strong and weak interactions of hyperons. Theoretical efforts have been devoted to understanding hypernuclear structure and production and decay mechanisms. There the most important ingredient is the hyperon-nucleon interactions. Several realistic potential models are in use widely\textsuperscript{[5–7]}, but their foundations are not solid and furthermore, there are some discrepancies among the models. For instance, the strengths of the spin-spin interaction vary significantly among the models. It seems urgent to establish the quantitative description of the hyperon-nucleon (YN) interactions.

In studying the YN interactions, it is natural to follow the description of nuclear force. The long-range part of the nuclear force is explained very well in terms of one-pion exchange mechanism, while heavy mesons as well as multi-pion exchanges are necessary for

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the medium range part of the nuclear force. One-boson exchange potential models, in which two- (and multi-) pion exchanges are taken into account as the $\sigma$ and $\rho$ exchanges, are fairly successful in accounting the large amount of data for nucleon-nucleon scattering \cite{8}. Yet the short-range part of the nuclear force is not fully treated, that is, repulsive cores (hard or soft) are introduced phenomenologically to explain the NN scattering phase shifts around $E \approx 100 - 200$ MeV in the center of mass system. Indeed, this is the region where the internal quark-gluon structure of the nucleon must be considered explicitly.

It was pointed out that the quark exchange force between two nucleons give significant repulsion at the short distance. The exchange force is induced by the quark antisymmetrization and therefore is nonlocal and of short-range determined by the size of the quark content of the nucleon. The most important feature of the quark exchange force is its dependence on the spin-flavor symmetry of two-baryon states. A close analogy is found in the hydrogen molecule, where two electrons orbit around two protons. As the total spin of the electrons specifies the symmetry of the spin wave function, the sign of the exchange force is determined according to the spin. The symmetric orbital state is allowed only for $S = 0$, while the exchange force is strongly repulsive for $S = 1$. Similar state dependencies appear in the quark exchange force, where the spin-flavor SU(6) symmetry determines the properties of the exchange interactions.

We applied the quark cluster model description of the short-range YN interactions\cite{9}. We found that the flavor singlet combination of $\Lambda\Lambda - N\Xi - \Sigma\Sigma$ has no repulsion induced by the quark exchange at the short distance. As this state is known to be favored by the magnetic part of the one-gluon exchange interaction (color-magnetic interaction), a bound or a resonance state called H dibaryon may exist\cite{11,12}. On the other hand, most other channels have strong repulsion at short distances.

Another interesting observation made in ref.\cite{9} is that the $S$ wave $\Sigma N$ interaction depends strongly on the total spin and isospin. The $\Sigma N$ ($I = 1/2, S = 0$) and $\Sigma N$ ($I = 3/2, S = 1$) states belong mainly to the $[51]$ irreducible representation of the spin-flavor SU(6) symmetry. This is the representation in which a Pauli forbidden state appears in the $L = 0$ orbital motion. The Pauli principle forbids two baryons to get together and thus gives a strong repulsion. The other spin-isospin states do not belong to this symmetry and therefore the short range repulsion is weaker. This qualitative argument was confirmed in realistic quark cluster model calculation of the YN interactions\cite{9,10}. Recent analyses by Niigata-Kyoto group show that the strong repulsion remains after combining the quark exchange interaction with the long-range meson exchange attraction\cite{10}. No experimental evidence is yet available to be able to confirm the strong state dependencies. More $\Sigma N$ scattering data are anticipated very much.

### 3. Antisymmetric Spin-Orbit Forces\footnote{This part of the work has been done in collaboration with Yoshihiro Tani\cite{2}.}

One of the interesting features of the hyperon-nucleon interactions is the properties of the spin-orbit force. The Galilei invariant spin-orbit force consists of symmetric LS (SLS) and antisymmetric LS (ALS) terms,

$$V_{SO} = V_{SLS} (\sigma_1 + \sigma_2) \cdot L + V_{ALS} (\sigma_1 - \sigma_2) \cdot L$$
Because the ALS operator \((\sigma_1 - \sigma_2) \cdot L\) is antisymmetric with respect to the exchange of two baryons, \(V_{ALS}\) should be zero between like baryons. In the nuclear force, ALS between proton and neutron breaks the isospin symmetry and is classified as a type IV charge symmetry breaking (CSB) force. Evidence of such a CSB force was given by measuring the difference between the proton and neutron analyzing powers in the \(n - p\) scattering experiments\(^{[13]}\). The results show that this force is very weak, supporting the isospin invariance of the nuclear force.

On the contrary, the ALS forces in the hyperon-nucleon interactions do not vanish even in the SU(3) flavor symmetric limit. For hyperon-nucleon systems, ALS seems as strong as the symmetric LS part. If the magnitudes of SLS and ALS are comparable, then the single particle LS force for one of the baryons, (ex. \(\Lambda\)) inside (hyper)nuclei is much smaller than that for the other baryon (nucleon). Since recent experiment suggests that the single particle LS force for \(\Lambda\) might be sizable contrary to the wide belief of vanishing LS force for \(\Lambda\), it is extremely important to pin down the magnitude of the two-body LS force. Here we study the properties of the YN ALS forces from the SU(3) symmetry point of view.

3.1. SU(3) Invariance

For the octet baryons, the baryon-baryon interactions can be classified in terms of the SU(3) irreducible representations given by

\[
8 \times 8 = 1 + 8_s + 27 + 10 + 10^* + 8_a
\]  

Among these six irreducible representations, first three, \(1, 8_s, \) and \(27,\) are symmetric under the exchange of two baryons and the other three \(10, 10^*\) and \(8_a,\) are antisymmetric\(^{[6]}\). Noting that two-baryon states are to be antisymmetric, and that the color wave function for the color-singlet baryons is always symmetric, we find that the symmetric (antisymmetric) flavor representations are combined only to antisymmetric (symmetric) spin-orbital states.

In baryon-baryon scattering, the ALS force induces the transition (mixing) between the spin singlet states \((^1P_1, ^1D_2, ^1F_3, \ldots)\) and the spin triplet states with the same \(L\) and \(J\) \((^3P_1, ^3D_2, ^3F_3, \ldots).\) The flavor symmetries of \(^1P_1\) state must be antisymmetric, \(10, 10^*\) and \(8_a,\) while that of \(^3P_1\) state is symmetric, \(1, 8_s\) or \(27.\)

If one assumes the SU(3) invariance of the strong interaction, different irreducible representations are not mixed. Therefore the only possible combination of symmetric and antisymmetric representations is \(8_s - 8_a.\) We conclude that the ALS in the SU(3) limit should only connect \(8_s\) and \(8_a.\)

The symmetry structure becomes clearer by decomposing the \(P\)-wave \(\Lambda N - \Sigma N\) \((I = 1/2),\) as a concrete example, into the SU(3) irreducible representations. The flavor symmetric states read

\[
\begin{pmatrix}
 \Lambda N \\
 \Sigma N
\end{pmatrix} (^3P_1) = \begin{pmatrix}
 \sqrt{\frac{9}{10}} & -\sqrt{\frac{1}{10}} \\
 -\sqrt{\frac{1}{10}} & -\sqrt{\frac{9}{10}}
\end{pmatrix} \begin{pmatrix}
 27 \\
 8_s
\end{pmatrix}
\]
In the SU(3) limit, the only surviving matrix element is \( \langle 8_a | P_1 | 8^*_s \rangle \). When we turn to the YN particle basis, we obtain the following relations in the SU(3) limit.

\[
\langle \Lambda N^1 P_1 | V | \Lambda N^3 P_1 \rangle = -\langle \Sigma N^{(1/2)} | P_1 | \Lambda N^3 P_1 \rangle \\
\langle \Lambda N^1 P_1 | V | \Sigma N^{(1/2)} 3 P_1 \rangle = -\langle \Sigma N^{(1/2)} | P_1 | \Sigma N^{(1/2)} 3 P_1 \rangle = 3 \langle \Lambda N^1 P_1 | V | \Lambda N^3 \rangle
\]

On the contrary, the \( \Sigma N (I = 3/2) \) system belongs purely to the SU(3) 27 and therefore the ALS matrix element vanishes in the SU(3) limit.

\[
\langle \Sigma N^{(3/2)} 1 P_1 | V | \Sigma N^{(3/2)} 3 P_1 \rangle = 0
\]

These relations come only from the SU(3) symmetry and is general for any ALS interactions regardless their origin. It is extremely interesting to note that (1) the ALS for \( \Sigma N^{(1/2)} \) is much stronger than and has different sign from that for \( \Lambda - N \), (2) the coupling of \( \Lambda N - \Sigma N^{(1/2)} \) is also strong, and (3) the ALS for \( \Sigma N \) depends strongly on the isospin or the charge states.

### 3.2. Quark Cluster Model Potential

The quark substructure of the baryon must play significant roles at short distances in the baryonic interactions\[14\]. The quark model symmetry is especially simple when two baryons sit on top of each other. For \( L = 1 \) states, the six-quark orbital state takes [51] symmetry, which can couple only to the [42] SU(6) representation due to the Pauli principle. Namely the [6] symmetric SU(6) states are not allowed to couple to [51] orbital configuration. The \( I = 1/2 \) \( \Lambda N - \Sigma N \) states at \( R \to 0 \) are given in the SU(6)/SU(3) symmetry basis by

\[
| S = 0 \rangle = \langle [42] 8_a \rangle + \langle [42] 10^* \rangle \\
| S = 1 \rangle = \langle [42] 8_s \rangle + \langle [42] 27 \rangle
\]

Again the ALS in the SU(3) limit can connect only the octet states, \( [42] 8_a \) and \( [42] 8_s \), and the ratio of the ALS interaction in the particle basis, \( \Lambda N - \Sigma N \) should follow eq.(3).

The quark-quark potential due to one-gluon exchange gives a \( q - q \) spin-orbit interaction. Its contribution to the YN ALS forces are calculated in the above mentioned limit \( R = 0 \) [13]. They are compared with the corresponding potential of the ordinary spin orbit force between the \( 3 P_1 \) states in Table [I].

The results show that the ALS forces due to the quark exchange force are as strong as the ordinary LS force of the same origin. Especially, the \( \Sigma N (I = 1/2) \) feels a stronger ALS force between \( S = 0 \) and \( S = 1 \), than the ordinary LS force between \( S = 1 \) states. This is a very interesting result. The quark exchange ALS force might play dominant role in the YN forces, as was already suggested and demonstrated by Kyoto-Niigata group [11].

Table [I] also shows the potential values when the SU(3) symmetry is broken by the mass difference of the strange quark and the \( ud \) quarks. The effects of the symmetry breaking are not so large that the results are essentially the same. Thus the above SU(3) relations of the ALS matrix elements remain valid qualitatively, \( i.e., \langle \Lambda N^1 P_1 | V | \Sigma N^{(1/2)} 3 P_1 \rangle \) and \( \langle \Sigma N^{(1/2)} 1 P_1 | V | \Sigma N^{(1/2)} 3 P_1 \rangle \) are large, while \( \langle \Sigma N^{(3/2)} 1 P_1 | V | \Sigma N^{(3/2)} 3 P_1 \rangle \) almost vanishes.
Table 1
ALS and LS matrix elements at $R = 0$ normalized by the overlapping matrix element.

| $(I = 1/2)$ | $\langle 1P_1|V_{ALS}|1P_1 \rangle$ (MeV) | $\langle 3P_1|V_{LS}|3P_1 \rangle$ (MeV) |
|-------------|---------------------------------|---------------------------------|
| $\Lambda N \leftrightarrow \Lambda N$ | 37 | 32 |
| $\Lambda N \leftrightarrow \Sigma N$ | 88 | 77 |
| $\Sigma N \leftrightarrow \Lambda N$ | -37 | -29 |
| $\Sigma N \leftrightarrow \Sigma N$ | -88 | -79 |

| $(I = 3/2)$ | $\langle 3P_1|V_{ALS}|3P_1 \rangle$ (MeV) | $\langle 3P_1|V_{LS}|3P_1 \rangle$ (MeV) |
|-------------|---------------------------------|---------------------------------|
| $\Sigma N \leftrightarrow \Sigma N$ | 0 | 1 |

3.3. Meson Exchange Potential

For meson exchange interactions exchanged mesons considered are either in the flavor singlet or octet representation (because they are $q\bar{q}$ states). The SU(3) factor for the (meson $M^a$)–(baryon $B_i$)–(baryon $B_j$) coupling, $T^a_{ij}$, has three choices, $\delta_{ij}$ for the flavor singlet meson ($a = 0$) and $F_{aij}$ or $D_{aij}$ for octet mesons ($a = 1$–8), where $F$ and $D$ are symmetric and antisymmetric SU(3) structure constants, respectively. Then the SU(3) invariant potential is proportional to

$$\sum_{a=1}^8 (T^a_{ij} \cdot T^a_{lm} + \text{exchange term})$$

Thus possible antisymmetric coupling is of the form $(F_{aij} \cdot D_{alm} - D_{aij} \cdot F_{alm})$. This term, however, vanishes because the ratio of the $f$ and $d$ couplings is fixed for each meson without depending on the choice of baryons $(ijlm)$. One exception is for the vector and the tensor couplings in the vector meson exchange force. According to the vector meson dominance, $F/D$ ratio for the vector and the tensor couplings are in general different and then terms like $(g_1 f_2 - f_1 g_2)(\sigma_1 - \sigma_2) \cdot L$ will survive, where $g_k$ ($f_k$) is the vector (tensor) coupling constant of a vector meson to a baryon $k$ ($k = 1$ or 2).

One can confirm the above symmetry consideration by a look on the ALS potential term in the Nijmegen potential [4], for instance. There the SU(3) symmetry is broken by differences among the baryon masses that cause factors like $(M^2_B - M^2_N)/4M_YM_N$. The equal baryon mass $M_Y = M_N$ in the SU(3) limit kills most terms, leaving

$$V_{ALS}(r) = \frac{g_Yg_N}{4\pi} \frac{m^3}{2M^2} \xi(x) \left[ \left( \frac{f}{g} \right)_N - \left( \frac{f}{g} \right)_Y \right] (\sigma_1 - \sigma_2) \cdot L$$

(8)

where $m$ is the meson mass, $x = nr$ and $\xi(x)$ is a radial function defined by

$$\xi(x) = \left( \frac{1}{x} + \frac{1}{x^2} \right) e^{-x} \frac{x}{x} \cdot$$

(9)

In the SU(3) limit, the coefficients of ALS force due to exchanges of the vector mesons, $\rho$, $\omega_8$ and $K^*$ are given for the SU(3) basis, and also for the particle basis in Table [4]. The common factor $\frac{1}{\sqrt{5}} f g (\alpha - \beta)$ contains $\alpha = F_g/(F_g + D_g)$ for the vector($g$)-coupling
and $\beta = F_f / (F_f + D_f)$ for the tensor(f)-coupling. All the matrix elements vanish when $\alpha = \beta$. Thus it is critical for ALS to have different $F/D$ ratios for the vector and tensor couplings.

Table 2

|            | $\rho$ | $\omega_8$ | $K^*$ | Total |
|------------|--------|------------|-------|-------|
| $\langle 8^a_1 P_1 | V_{ALS} | 8^a_1 3P_1 \rangle$ | 8 | 4 | 8 | 20 |
| $\langle 8^a_1 P_1 | V_{ALS} | 27^3 P_1 \rangle$ | 6 | -2 | -4 | 0 |
| $\langle 10^* 1 P_1 | V_{ALS} | 8^a_1 3P_1 \rangle$ | -2 | -2 | 4 | 0 |
| $\langle 10^* 1 P_1 | V_{ALS} | 27^3 P_1 \rangle$ | -4 | -4 | 8 | 0 |
| $\langle \Lambda N^{(1/2)} 1 P_1 | V_{ALS} | \Lambda N^{(1/2)} 3P_1 \rangle$ | 0 | -1 | 0 | -1 |
| $\langle \Sigma N^{(1/2)} 1 P_1 | V_{ALS} | \Sigma N^{(1/2)} 3P_1 \rangle$ | 2 | 1 | 0 | 3 |
| $\langle \Sigma N^{(1/2)} 1 P_1 | V_{ALS} | \Lambda N^{(1/2)} 3P_1 \rangle$ | -1 | 0 | 2 | 1 |
| $\langle \Lambda N^{(1/2)} 1 P_1 | V_{ALS} | \Sigma N^{(1/2)} 3P_1 \rangle$ | -1 | 0 | -2 | -3 |
| $\langle 10^* 1 P_1 | V_{ALS} | 27^3 P_1 \rangle$ | 1 | -1 | 0 | 0 |
| $\langle \Sigma N^{(3/2)} 1 P_1 | V_{ALS} | \Sigma N^{(3/2)} 3P_1 \rangle$ | 1 | -1 | 0 | 0 |

One sees that except for the $8^a - 8^a$ matrix element, the sum of the $\rho$, $\omega_8$ and $K^*$ contributions vanishes as expected from the SU(3) symmetry. It is also interesting to note that the difference between the $\Lambda N \rightarrow \Lambda N$ and $\Sigma N \rightarrow \Sigma N$ ALS matrix elements come from the isovector $\rho$ exchange and also that the difference between $\Lambda N \rightarrow \Sigma N$ and $\Sigma N \rightarrow \Lambda N$ is caused by the sign change of $K^*$ exchange. For $\langle \Lambda N^{1} P_{1} | V_{ALS} | \Sigma N^{(1/2)} 3P_{1} \rangle$, the $\rho$ and $K^*$ exchanges are added up while they tend to cancel for $\langle \Sigma N^{(1/2)} 1 P_{1} | V_{ALS} | \Lambda N^{3} P_{1} \rangle$.

### 3.4. SU(3) Breaking

When the SU(3) symmetry breaking is taken into account, the coefficients in Table 2 for $\rho$, $\omega_8$ and $K^*$ cannot simply be summed up, as the interaction range for $K^*$ is significantly shorter than that for $\rho$ or $\omega_8$. Furthermore the differences among the octet baryon masses generate various terms that would vanish in the SU(3) limit. The full form of the ALS potential for the scalar, pseudoscalar and vector exchanges are given in the following.

$$V_{PS}^{ALS}(r) = g_{13}g_{24} \frac{m^3}{4\pi} \xi(x) \frac{1}{8} \left( \frac{1}{M_1 M_4} - \frac{1}{M_2 M_3} \right) \mathbf{L} \cdot (\mathbf{\sigma}_1 - \mathbf{\sigma}_2) P_\sigma$$

$$V_{S}^{ALS}(r) = - g_{13}g_{24} \frac{m^3}{4\pi} \xi(x) \frac{1}{8} \left( \frac{1}{M_1 M_3} - \frac{1}{M_2 M_4} \right) \mathbf{L} \cdot (\mathbf{\sigma}_1 - \mathbf{\sigma}_2)$$
\[ V_{V^{ALS}}(r) = -\frac{m^3}{4\pi} \xi(x) \frac{1}{2} \mathbf{L} \cdot (\mathbf{\sigma}_1 - \mathbf{\sigma}_2) \]

\[
= \left[ (g_{13} + \frac{M_1 + M_3}{2M} f_{13}) (g_{24} + \frac{M_2 + M_4}{2M} f_{24}) \right] \frac{1}{4} \left( \frac{1}{M_2 M_3} - \frac{1}{M_1 M_4} \right) \]

\[
\times \left( \frac{1}{M_1} + \frac{1}{M_2} + \frac{1}{M_3} + \frac{1}{M_4} \right) \left( \frac{1}{M_1 M_3} - \frac{1}{M_2 M_4} \right) \]

\[
+ \frac{g_{13} f_{24} - f_{13} g_{24}}{4M} \left( \frac{1}{M_1} + \frac{1}{M_2} + \frac{1}{M_3} + \frac{1}{M_4} \right) \left( 1 + \frac{(M_3 - M_1)(M_4 - M_2)}{m^2} \right) \]

\[
+ \frac{g_{13} f_{24}}{4M^2} \left( \frac{1}{2} \left( \frac{M_2^2 + M_4^2}{M_2 M_4} - \frac{M_1^2 + M_3^2}{M_1 M_3} \right) \right) \]

\[
- \frac{m^2}{16} \left( \frac{1}{M_1 M_3} - \frac{1}{M_2 M_4} \right) \]

\[
+ \frac{1}{32} \left( \frac{1}{M_1^2} + \frac{1}{M_3^2} - \frac{1}{M_2^2} - \frac{1}{M_4^2} \right) \left( \frac{M_2^2 + M_4^2 - M_1^2 - M_3^2}{M_1 M_2 M_3 M_4} \right) \]

\[
\frac{1}{2} \left( \mathbf{\sigma}_1 \cdot \mathbf{\sigma}_2 \right) \]

where the scattering of \( M_1 + M_2 \rightarrow M_3 + M_4 \) is considered, while \( M \) is the mass of the proton. The terms that contains the spin exchange operator,

\[ P_\sigma = \frac{1 + \mathbf{\sigma}_1 \cdot \mathbf{\sigma}_2}{2} \]

come from the exchange part in the \( K \) and \( K^* \) exchanges.

Fig. 1 shows the ALS part of the one-boson exchange potential (OBE) and the adiabatic potential in the quark cluster model (QCM) for the \( I = 3/2 \) \( \Sigma N \) channel, i.e., \( \langle \Sigma N^{(3/2)} \ 1P_1 | V_{ALS} | \Sigma N^{(3/2)} \ 3P_1 \rangle \). It can be compared with the corresponding ordinary L S (SLS) part for \( \langle \Sigma N^{(3/2)} \ 3P_1 | V_{SLS} | \Sigma N^{(3/2)} \ 3P_1 \rangle \), shown in Fig. 2. The curves in these figures show various cases in SU(3) breaking. The potentials in the SU(3) symmetric limit are labeled by “Bs-Ms” for OBE and “QCM-s” for QCM. The ALS potentials are zero in this limit. The symmetry is broken in “QCM-d” by the quark mass difference, while in “Bd-Md” both the meson mass differences and the baryon mass differences are taken into account. It is clear that the ALS is strongly suppressed even when the SU(3) symmetry breaking is considered. SU(3) breaking effect is small also in SLS. In general, we observe that the ALS potential is very weak for the two-baryon channels in which the ALS vanishes in the SU(3) limit. In these figures, potential parameters are chosen from the Nijmegen model D potential in OBE and ref. [13] for QCM.

Figs. 3 and 4 shows the OBE and QCM adiabatic potentials in the \( \Lambda N^{(1P_1)} \rightarrow \Sigma N^{(1/2)} \ 3P_1 \) (Fig. 3), and \( \Sigma N^{(1/2)} \ 1P_1 \rightarrow \Lambda N^{(3P_1)} \) (Fig. 4). They show that the ALS potentials are rather strong. It is also indicated that the QCM gives stronger ALS than OBE and also that the effects of SU(3) breaking is less for QCM.

4. Conclusion

We have studied the SU(3) symmetry of the ALS interactions in the YN force, and found that the SU(3) symmetry is rather good in accounting the properties of ALS forces in
Figure 1. The OBE and QCM potentials in $\langle \Sigma N^{(3/2)} 1 P_1 | V_{ALS} | \Sigma N^{(3/2)} 3 P_1 \rangle$.

Figure 2. The OBE and QCM potentials in $\langle \Sigma N^{(3/2)} 3 P_1 | V_{SLS} | \Sigma N^{(3/2)} 3 P_1 \rangle$.

Figure 3. The OBE and QCM potentials in $\langle \Lambda N^{(1/2)} 1 P_1 | V_{ALS} | \Sigma N^{(1/2)} 3 P_1 \rangle$.

Figure 4. The OBE and QCM potentials in $\langle \Sigma N^{(1/2)} 1 P_1 | V_{ALS} | \Lambda N^{(1/2)} 3 P_1 \rangle$. 
various baryon channels. The YN ALS forces due to the quark exchange are significantly large, comparable to the ordinary LS force of the same origin. The meson exchange force almost vanishes except for a term proportional to the difference in $F/D$ ratios of the vector and tensor couplings of the vector mesons. Thus the ALS interaction has a shorter range than the ordinary LS force, or the tensor forces. It is extremely interesting to pin down the strengths and the properties of the ALS forces in the YN sector in determining the origin of the baryonic forces. Further theoretical and experimental studies of the YN spin-orbit interactions are very much encouraged.

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