Roles of Quarks in Strong and Weak YN Interactions

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The short range parts of the strong and weak hyperon-nucleon interactions are studied with quark substructure of the baryons taken into account. A summary of the quark cluster model calculations of the baryon-baryon interactions is presented. It is pointed out that the spin-flavor symmetry structure determines the qualitative behaviors of the short-range YN interactions. We also show that the spin-orbit YN forces give characteristic behaviors, which are distinct from the meson exchange forces. Weak decays of $\Lambda$ in hypernuclei are described by the direct quark mechanism. We find that the neutron induced nonmesonic decay is enhanced due to the direct quark transition. Importance of the $\Delta I = 3/2$ decay amplitudes is emphasized both for the nonmesonic weak decays and the $\pi^+$ emission in the $\Lambda$ decay.

1 Introduction

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. The hyperon brings a new flavor, strangeness, to the up-down world of nuclei. Several interesting questions, such as the role of the Pauli principle for the hyperon, possible existence of dibaryon resonances, etc., are yet to be answered. It is therefore desirable to understand the behaviors of the hyperon in nuclear environment in the context of QCD. We here consider the quark substructure of the baryons and study the roles of the explicit quark degrees of freedom. The spin and flavor symmetry of quarks plays the most important role as the SU(3) symmetry was the driving force for developing the quark model of the hadrons. We mostly concentrate on general arguments that are based only on the symmetry structure, and for a numerical calculation we introduce a model, where the simple SU(3) constituent quark model and one-gluon exchange interactions as well as the confining force among the quarks are employed.

The weak interaction is another interesting subject, which can be studied in detail in the strangeness system. Hypernuclei are rich sources of weak strangeness decay in nuclear medium, where a new decay mode, $\Lambda N \rightarrow N N$
is realized. We point out that the $\Delta I = 1/2$ rule is not necessarily satisfied for nonmesonic decays of hyperons, and that the neutron induced nonmesonic decay is strongly enhanced due to the direct quark transition. It is also found that $\pi^+$ decay of hypernuclei is generated mainly by the $\Delta I = 3/2$ weak transitions.

In sect. 2, we present the quark model description of the baryon-baryon interaction including hyperons. The antisymmetric spin-orbit force is examined in sect. 3 from the SU(3) flavor symmetry point of view. In sect. 4, we introduce the direct quark mechanism for nonmesonic weak decays of $\Lambda$. In sect. 5, the $\pi^+$ decay of hypernuclei and the role of the $\Delta I = 3/2$ weak transitions are discussed.

2 Short range part of YN Interactions

We first summarize the main results of the quark model study of the hyperon-nucleon (YN) interactions. Although the details depend on models of the quark dynamics we choose, some robust qualitative features are derived from the symmetry argument.

In studying the YN interactions, it is natural to follow descriptions of the 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 the medium range part. One-boson exchange potential models, in which two- (and multi-) pion exchanges are taken into account in terms of $\sigma$ and $\rho$ exchanges, are fairly successful in accounting the large amount of data for nucleon-nucleon scattering. Yet the short-range part of the nuclear force is not fully understood microscopically, that is, a repulsive core (hard or soft) is often 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 play important roles.

It was pointed out that the quark exchange force between two nucleons gives strong repulsion at short distances. 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 to the short-range YN interactions. We found that the flavor singlet combination of $\Lambda\Lambda - N\Xi - \Sigma\Sigma$ has no repulsion induced by the quark exchange. 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. On the other hand, most other YN and YY channels have strong repulsion at short distances.

Another interesting observation made in ref. is that the $S$ wave $\Sigma N$ interaction depends strongly on the total spin and isospin. The $S$-wave $\Sigma N$ ($I = 1/2, S = 0$) and $\Sigma N$ ($I = 3/2, S = 1$) states belong mainly to the $[51]$ symmetric 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 closer 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 calculations of the YN interactions. 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. 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

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,

$$V_{SO} = V_{SLS}(\sigma_1 + \sigma_2) \cdot L + V_{ALS}(\sigma_1 - \sigma_2) \cdot L$$

$$= (V_{SLS} + V_{ALS}) \sigma_1 \cdot L + (V_{SLS} - V_{ALS}) \sigma_2 \cdot L \quad (1)$$

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. 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 SLS part. If the magnitudes of SLS and ALS are comparable, then the single particle LS force for one of the baryons, (ex. Λ) inside (hyper)nuclei is much weaker than that for the other baryon (nucleon). Since recent experiment suggests that the single particle LS force for Λ might be sizable contrary to the wide belief of vanishing LS force for Λ, it is extremely important to pin down the magnitude of the two-body LS force.

3.1 SU(3) symmetry for ALS

First we study the properties of the YN ALS forces from the SU(3) symmetry point of view. 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. 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

\[
\begin{pmatrix}
\Lambda N \\
\Sigma N
\end{pmatrix} \left(^3P_1\right) = \begin{pmatrix}
\sqrt{9/10} & -\sqrt{1/10} \\
-\sqrt{1/10} & -\sqrt{9/10}
\end{pmatrix} \begin{pmatrix}
27 \\
8_s
\end{pmatrix}
\]  

while the antisymmetric ones are

\[
\begin{pmatrix}
\Lambda N \\
\Sigma N
\end{pmatrix} \left(^1P_1\right) = \begin{pmatrix}
-\sqrt{1/2} & -\sqrt{1/2} \\
-\sqrt{1/2} & \sqrt{1/2}
\end{pmatrix} \begin{pmatrix}
10^* \\
8_a
\end{pmatrix}
\]
In the SU(3) limit, the only surviving matrix element is \( \langle 8_a^1 P_1 | V | 8_s^3 P_1 \rangle \). When we turn to the YN particle basis, we obtain the following relations in the SU(3) limit.

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

(5)

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.

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

(6)

These relations come only from the SU(3) symmetry and is general for any ALS interactions regardless their origin. We especially 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 approach

In the quark cluster model approach, the one-gluon exchange gives \( q - q \) spin-orbit interaction, and its contribution to the YN ALS forces is evaluated easily in the adiabatic approximation.23 The results for the \( P \) wave states at \( R = 0 \), where two baryons sit on top of each other, are compared with the corresponding symmetric spin orbit (SLS) forces in Table 1. The results show that the ALS forces due to the quark exchange are as strong as the SLS 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 SLS force between \( S = 1 \) states. On the other hand, the ALS force vanishes in the \( \Sigma N (I = 3/2) \), such as \( \Sigma^+ p \) system.

Table 1 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. (See ref.2 for details.) This interesting result, that the quark exchange ALS force will play dominant role in the YN forces, has been suggested and demonstrated by Kyoto-Niigata group in a realistic model of the YN interaction based on the quark cluster model.24

### 3.3 Meson exchange force

The ALS YN force in the meson exchange potential is known to be weak compared to SLS. The reason can again be understood qualitatively from the
Table 1: ALS and LS matrix elements at \( R = 0 \) normalized by the overlapping matrix element.

| Coupling              | \( \langle 1P_1|V_{ALS}|3P_1 \rangle \) (MeV) | \( \langle 3P_1|V_{SLS}|3P_1 \rangle \) (MeV) |
|-----------------------|---------------------------------------------|---------------------------------------------|
| \( I = 1/2 \)         |                                             |                                             |
| \( \Lambda N \leftrightarrow \Lambda N \) | 37 | 32 | \( -74 \) | \( -55 \) |
| \( \Lambda N \leftrightarrow \Sigma N \) | 88 | 77 | 33 | 29 |
| \( \Sigma N \leftrightarrow \Lambda N \) | \( -37 \) | \( -29 \) | 33 | 29 |
| \( \Sigma N \leftrightarrow \Sigma N \) | \( -88 \) | \( -79 \) | 22 | 22 |
| \( I = 3/2 \)         |                                             |                                             |
| \( \Sigma N \leftrightarrow \Sigma N \) | 0 | 1 | \( -95 \) | \( -94 \) |

SU(3) symmetry. There the exchanged mesons are either in the flavor singlet or octet representation (because they are \( q \bar{q} \) states). The SU(3) factor for the \( (\text{meson } M_a) \rightarrow (\text{baryon } B_i) \rightarrow (\text{baryon } B_j) \) coupling, \( T_{ij}^a \), has three choices, \( \delta_{ij} \) for the flavor singlet meson \((a = 0)\) and \( F_{aij} \) or \( D_{aij} \) for octet mesons \((a = 1 \cdots 8)\), where \( F \) and \( D \) are symmetric and antisymmetric SU(3) structure constants, respectively. Then the SU(3) invariant potential is proportional to

\[
\sum_{a=0}^{8} (T_{ij}^a \cdot T_{lm}^a + \text{exchange term})
\]

One sees that the only 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 \( (g_1 f_2 - f_1 g_2)(\boldsymbol{\sigma}_1 - \boldsymbol{\sigma}_2) \cdot \mathbf{L} \) will survive, where \( g_k \) \((f_k)\) is the vector \( (\text{tensor}) \) coupling constant of a vector meson to a baryon \( k \) \((k = 1 \text{ or } 2)\). Thus the contributions of the meson exchanges to ALS YN force are very limited in the SU(3) limit, and we find that they are much weaker than the SLS counterparts. The situation again does not change even if the SU(3) symmetry is broken. See ref. [4] for details.

Thus it is important to pin down the strengths and the properties of the ALS forces in the YN sector as they indicate the origin of the short-distance baryonic forces. Further studies of the YN spin-orbit interactions are very much encouraged.
4 Direct Quark Process for the Weak Decay of Λ

Recent experimental and theoretical studies of weak decays of hypernuclei have generated renewed interest on nonleptonic weak interactions of hadrons. A long standing problem is the dominance of $\Delta I = 1/2$ amplitudes in the strangeness changing transitions. The decays of kaons, and $\Lambda$, $\Sigma$ hyperons are dominated by the $\Delta I = 1/2$ transition but it is not clear whether this dominance is a general property of all nonleptonic weak interactions. In fact, the weak effective interaction which is derived from the standard model including the perturbative QCD corrections contains a significantly large $\Delta I = 3/2$ component. It is therefore believed that nonperturbative QCD corrections, such as hadron structures and reaction mechanism, are responsible for suppression of $\Delta I = 3/2$, and/or enhancement of $\Delta I = 1/2$ transition amplitudes.

From this viewpoint, decays of hyperons inside nuclear medium provide us with a unique opportunity to study new types of nonleptonic weak interaction, that is, two- (or multi-) baryon processes, such as $\Lambda N \rightarrow NN$, $\Sigma N \rightarrow NN$, etc. These transitions constitute the main branch of hypernuclear weak decays because the pionic decay $\Lambda \rightarrow N\pi$ is suppressed due to the Pauli exclusion principle for the produced nucleon.

A conventional picture of the two-baryon decay process, $\Lambda N \rightarrow NN$, is the one-pion exchange between the baryons, where $\Lambda N\pi$ vertex is induced by the weak interaction. In $\Lambda N \rightarrow NN$, the relative momentum of the final state nucleon is about 400 MeV/c, much higher than the nuclear Fermi momentum. The nucleon-nucleon interaction at this momentum is dominated by the short-range repulsion due to heavy meson exchanges and/or to quark exchanges between the nucleons. It is therefore expected that the short-distance interactions will contribute to the two-body weak decay as well. Exchanges of $K$, $\rho$, $\omega$, and $K^*$ mesons and also correlated two pions in the nonmesonic weak decays of hypernuclei have been studied, and it is found that the kaon exchange is significant, while the other mesons contribute less.

Several studies have been made on effects of quark substructure. In our recent analyses, we employ an effective weak hamiltonian for quarks, which takes into account one-loop perturbative QCD corrections to the $W$ exchange diagram in the standard model. We evaluated the effective hamiltonian in the six-quark wave functions of the two baryon systems and derived the “direct quark (DQ)” weak transition potential for $\Lambda N \rightarrow NN$. Our analysis showed that the DQ contribution is significantly large compared to the conventional pion exchange amplitudes, and shows some qualitatively distinct features. It largely improves the discrepancy between the meson-exchange theory and experimental data for the ratio of the neutron- and proton-induced
decay rates of light hypernuclei. It was also found that the $\Delta I = 3/2$ component of the effective Hamiltonian gives a sizable contribution to $J = 0$ transition amplitudes.

4.1 Nonmesonic Weak Decay of $\Lambda$ in Nuclei

The DQ transition takes place only when $\Lambda$ overlaps with a nucleon in hypernuclei and therefore predominantly in the relative $S$-states of $\Lambda N$ systems. The two-body transition potentials in all the possible channels with the initial $\Lambda N(L = 0)$ to the final $NN(L = 0, 1)$ states for the DQ mechanism are computed. Because of quark antisymmetrization effects, the DQ transition potentials contain a nonlocal component and as the transition may break the parity invariance, it also contains a derivative term. The general form of the transition potential is

$$V_{s's'JsJ}(r, r') = \left\langle NN : s's'J | V(r', r) | \Lambda N : sJ \right\rangle = V_{loc}(r) \frac{\delta(r - r')}{r^2} + V_{der}(r) \frac{\delta(r - r')}{r^2} \partial_r + V_{nonloc}(r', r) (7)$$

We compare the DQ potential with conventional meson exchange transition potentials, such as one-pion exchange (OPE). Because the OPE potential is determined phenomenologically, the relation between DQ and OPE is not trivial. In order to fix the relative phase of the two, we relate the $\pi N \Lambda$ coupling constant to a baryon matrix element of the weak Hamiltonian for quarks by using a soft-pion relation,

$$\lim_{q \to 0} \langle \pi^0(q) | H_{PV} | \Lambda \rangle = \frac{i}{f_\pi} \langle n | [Q^3_\pi, H_{PV}] | \Lambda \rangle = -\frac{i}{2f_\pi} \langle n | H_{PC} | \Lambda \rangle (8)$$

Here we use the relation, $[Q^3_\pi, H_W] = -[I^3, H_W]$, which is satisfied as the weak Hamiltonian $H_W$ consists only of left-hand currents and the flavor-singlet right-hand currents.

In ref. [3, 4], we calculated the nonmesonic decay rates of $\frac{2}{3}H_e$, $\frac{4}{3}H_e$, and $\frac{4}{3}H$ in the DQ and OPE mechanisms. The $S$-shell hypernuclei are most suitable for the study of the microscopic mechanism of the weak decay as their wave functions are relatively simple and contain only $\Lambda N(L = 0)$ states. They also enable us to select spin-isospin components for the weak decay. We found that DQ gives the major contribution in $J = 0$ transitions and therefore enhances the neutron induced decay rates. The superposition of the DQ and OPE shows a good agreement with available experimental data. We found that the $\Delta I = 3/2$ components in the $J = 0$ transitions are significant. We also pointed out that the nonmesonic decay rate of $\frac{4}{3}H$ is strongly enhanced by
DQ, and therefore its experimental data are critically important to confirm the DQ mechanism for the nonmesonic weak decay of Λ in nuclei.

We here present results of the calculation of the Λ decay in nuclear matter. We assume the $p-n$ symmetric nuclear matter with realistic short-range correlation of Λ and N. The results are summarized in Table 2. We compare the results of several different combinations of the meson exchanges and the direct quark processes.

Table 2: Nonmesonic decay rates of Λ in nuclear matter (in units of $\Gamma_\Lambda$). The form factors are taken into account for the meson exchanges, where the “hard” pion has $\Lambda_\pi = 1300\text{MeV}$ and the “soft” pion has $\Lambda_\pi = 800\text{MeV}$. The “all” includes $\pi, K, \eta, \rho, \omega,$ and $K^*$ meson exchanges.

|                | total  | $\Gamma_p$ | $\Gamma_n$ | $\Gamma_n/\Gamma_p$ | $P\text{V}/PC$ |
|----------------|--------|------------|------------|----------------------|----------------|
| $\pi$(hard)    | 2.575  | 2.354      | 0.221      | 0.094                | 0.337          |
| $\pi$(soft)    | 1.796  | 1.653      | 0.143      | 0.086                | 0.279          |
| $\pi$(hard)+K  | 1.099  | 1.076      | 0.024      | 0.022                | 0.631          |
| $\pi$(soft)+K  | 0.666  | 0.645      | 0.021      | 0.032                | 0.638          |
| $\pi$(hard)+all| 0.928  | 0.731      | 0.196      | 0.268                | 0.369          |
| $\pi$(soft)+all| 0.608  | 0.444      | 0.164      | 0.370                | 0.255          |
| DQ             | 0.418  | 0.202      | 0.216      | 1.071                | 6.759          |
| DQ+π(hard)     | 3.609  | 2.950      | 0.658      | 0.223                | 0.856          |
| DQ+π(soft)     | 2.661  | 2.147      | 0.514      | 0.239                | 0.902          |
| DQ+π(hard)+K   | 1.766  | 1.495      | 0.271      | 0.181                | 1.602          |
| DQ+π(soft)+K   | 1.164  | 0.962      | 0.202      | 0.210                | 2.000          |
| DQ+π(hard)+all | 1.507  | 1.123      | 0.384      | 0.342                | 1.471          |
| DQ+π(soft)+all | 1.020  | 0.734      | 0.286      | 0.390                | 1.584          |

We notice that the kaon exchange reduces the proton induced decay rates to more than factor two. This mainly comes from the suppression of the tensor transition, $\Lambda N :^3 S_1 \to NN :^3 D_1$. On the other hand, the neutron induced decays are much too small in the $\pi$ and $\pi+K$ exchanges, which is the main cause of the small $n/p$ ratio. The DQ mechanism, however, enhances the neutron induced decays and therefore improves the $n/p$ ratio. It is also shown that the parity violating decay is dominant in DQ transition, where the main component is the transition, $\Lambda N :^3 S_1 \to NN :^3 P_1$.

One sees in Table 2 that the pion exchange contribution strongly depends on the choice of the form factor. Some previous work uses a hard form factor with the cut off $\Lambda_\pi = 1300\text{MeV}$ according to the Bonn potential. But such form factor in general gives too large OPE contribution especially in the tensor...
transition, $\Lambda N ;^3 S_1 \rightarrow NN ;^3 D_1$. We here employ a softer form factor, $\Lambda \approx 800$ MeV, which is suggested by an analysis of the pion production in $NN$ scattering. It seems that the softer form factor is more appropriate to reproduce experimental values of the proton induced decay rates.

Contribution beyond $\pi$ and $K$ mesons, especially the vector mesons, contain some ambiguities. For instance, the values and even the signs of the weak coupling constants are not determined phenomenologically. They require SU(6) ansatz. They also depend strongly on the form factors, which are not well known. It is also questionable whether the DQ mechanism and the vector meson exchanges are independent and can be superposed. Here we follow the prescriptions given in ref. $^{20}$ for the vector exchange potentials and assume that there is no double counting in superposing DQ with the vector meson exchanges. We find a large $K^*$ contribution, which enhances the neutron induced decay rate and thus improves the $n/p$ ratio. We, however, do not think that this is the “final” result for the vector meson contribution because of the above mentioned ambiguities and unknown factors.

The reversed process, $pn \rightarrow \Lambda p$, which is the hyperon weak production in the $pn$ scattering, is also very interesting. We calculated the cross section of the $\Lambda$ production in the quark cluster model, taking the full six-quark wave function into account. The results will be published elsewhere.$^{24}$

5 \quad $\pi^+$ Decay of Hypernuclei

In this section, we study low energy $\pi^+$ emission in hypernuclear weak decays. We point out that the soft $\pi^+$ decay is directly related to $\Delta I = 3/2$ part of nonmesonic weak decays according to the soft pion theorem.

The $\pi^+$ emission from light hypernuclei, for instance, $^4\Lambda He$, has puzzled us for a long time. Rather old experimental data suggest that the ratio of $\pi^+$ and $\pi^-$ emission from $^4\Lambda He$ is about 5%. This small ratio is expected because the free $\Lambda$ decays only into $p\pi^-$ and $n\pi^0$. The $\pi^+$ emission requires an assistance of a proton, i.e., $\Lambda + p \rightarrow n + n + \pi^+$.

Several microscopic mechanisms for the $\pi^+$ emission have been considered in literatures.$^{26,27,28}$ The most natural one is $\Lambda \rightarrow n\pi^0$ decay followed by $\pi^0 p \rightarrow \pi^+ n$ charge exchange reaction. It was evaluated for realistic hypernuclear wave functions and found to explain only 1.2% for the $\pi^+ / \pi^-$ ratio.$^{26}$ Another possibility is to consider $\Sigma^+ \rightarrow \pi^+ n$ decay after the conversion $\Lambda p \rightarrow \Sigma^+ n$ by the strong interaction. It was found, however, that the free $\Sigma^+$ decay which is dominated by $P$-wave amplitude, gives at most 0.2% for the $\pi^+ / \pi^-$ ratio. Indeed it is clear that the $\Sigma^+$ mixing and its free decay is not the main mechanism, for experimental data suggest that the $\pi^+$ emission
is predominantly in the S-wave with the energy less than 15 MeV. Recently, it was proposed that a two-body process \( \Sigma^+ N \rightarrow nN\pi^+ \) must be important in the \( \frac{1}{2}He \) decay. But its microscopic mechanism is not specified.

Here, we would like to show that the \( \Delta I = \frac{3}{2} \) two-baryon transition amplitudes are directly related to the S-wave \( \pi^+ \) emission from hypernuclei. The relation of these two amplitudes is derived from the soft-pion theorem and is a result of the chiral structure of the weak interaction.

The soft-pion theorem for the process \( \Lambda p \rightarrow nn\pi^+ (q \rightarrow 0) \) gives a similar relation to the one we use in the previous section,

\[
\lim_{q \to 0} \langle nn\pi^+(q)|H_W|\Lambda p \rangle = \frac{i}{\sqrt{2}f_\pi} \langle nn|[I^-, H_W]|\Lambda p \rangle \tag{9}
\]

As \( H_W \) changes the third component of the isospin by \(-1/2\) when it converts \( \Lambda p \) (\( I_3 = +1/2 \)) to \( nn\pi^+ \) (\( I_3 = 0 \)), it may contain \( H_W(\Delta I = 1/2, \Delta I_z = -1/2) \) and \( H_W(\Delta I = 3/2, \Delta I_z = -1/2) \). Now it is easy to see that \( \Delta I = 1/2 \) part vanishes in eq.\( (9) \) as

\[
[I^-, H_W(\Delta I = 1/2, \Delta I_z = -1/2)] = 0 \tag{10}
\]

\[
[I^-, H_W(\Delta I = 3/2, \Delta I_z = -1/2)] = \sqrt{3} H_W(\Delta I = 3/2, \Delta I_z = -3/2) \tag{11}
\]

We then obtain

\[
\lim_{q \to 0} \langle nn\pi^+(q)|H_W|\Lambda p \rangle = \frac{i\sqrt{3}}{\sqrt{2}f_\pi} \langle nn|H_W(\Delta I = 3/2, \Delta I_z = -3/2)|\Lambda p \rangle \tag{12}
\]

Thus we conclude that the soft \( \pi^+ \) emission in the \( \Lambda \) decay in hypernuclei is caused only by the \( \Delta I = 3/2 \) component of the strangeness changing weak hamiltonian. In other words, the \( \pi^+ \) emission from hypernuclei probes the \( \Delta I = 3/2 \) transition of \( \Lambda N \rightarrow NN \).

Now we understand why the previous attempts to explaining the \( \pi^+/\pi^- \) ratio failed. Both the charge exchange process and the \( \Sigma^+ \) decay are induced by the \( \Delta I = 1/2 \) part of the hamiltonian and therefore cannot emit low-energy \( \pi^+ \). In fact, the reason why the S-wave \( \Sigma^+ \rightarrow n\pi^+ \) decay is very small is again that only the \( \Delta I = 3/2 \) amplitude can induce this decay for soft (S-wave) \( \pi^+ \).

In the same way, two-body \( \Sigma^+ \) decay will not contribute unless it is induced by a \( \Delta I = 3/2 \) weak interaction.

Supplying the \( \Lambda N \rightarrow NN \) two-body decay amplitudes given in the previous section, we now calculate the \( \pi^+ \) decay rates due to the two-body processes in the soft pion limit. Similar calculation can be done for the soft \( \pi^- \) decay in the two-body processes, where also the \( \Delta I = 1/2 \) components contribute.
Table 3: Soft $\pi^+$ and $\pi^-$ decay rates in arbitrary units.

|       | DQ $+\pi + K$ | DQ $+\pi$ |
|-------|---------------|-----------|
| $^4\Lambda$He $\rightarrow \pi^-$  | 193.9       | 208.5     |
|       | $\rightarrow \pi^+$ | 65.4     | 65.4     |
|       | $\pi^+/\pi^-$     | 34%       | 31%       |
| $^4\Lambda$H $\rightarrow \pi^-$  | 169.1       | 211.1     |
|       | $\rightarrow \pi^+$ | 130.6   | 130.6     |
|       | $\pi^+/\pi^-$     | 77%       | 62%       |
| $^5\Lambda$He $\rightarrow \pi^-$  | 185.5       | 216.9     |
|       | $\rightarrow \pi^+$ | 65.4     | 65.4     |
|       | $\pi^+/\pi^-$     | 35%       | 30%       |

Besides the direct quark process, the one-pion and one-kaon exchanges are included in the $\Delta I = 1/2$ amplitudes, while the $\Delta I = 3/2$ parts are purely from the direct quark mechanism. The results are shown in Table 3. We obtain the ratios of the $\pi^+$ emission to the $\pi^-$ emission from the two-body processes are as large as 77% in $^4\Lambda$H, and 30-35% for $^4\Lambda$He and $^5\Lambda$He.

In order to compare these ratios with experiment, we need to consider the threshold differences of these decays carefully, as the phase space volumes for the $\pi^+$ and $\pi^-$ decays are often largely different. We especially note, however, that the $^4\Lambda$He decay has a symmetric phase space volumes for $\pi^+$ and $\pi^-$ emissions and therefore the $\pi^+$ decay branch is most easily observed. In conclusion, it is extremely interesting to study the $\pi^+$ emission carefully so that the $\Delta I = 3/2$ component of the weak $\Lambda$ decay is confirmed.

6 Conclusion

We have exhibited several examples which represent the roles of explicit quark content of the baryons in strangeness nuclear physics. Both the strong and weak interactions in the YN systems show characteristic features of the quark substructure. It is in contrast to the NN system, where most phenomena can be accounted either with or without explicit quarks. We hope that the YN system can distinguish and enlighten the effects of quark substructure much more clearly. To this end, further efforts both in experimental and theoretical studies are necessary.

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