Strong and Weak Interactions of Strange Hadrons

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I review hadronic processes involving strange hadrons, especially hyperons from the quark structure point of view. The strong interaction of quarks expects several important new features when the strangeness is introduced upon the non-strange degrees of freedom. I concentrate on the instanton induced interaction, effects of the Pauli principle for strangeness and the nonleptonic weak interaction of strangeness. The single hadron as well as the two baryon systems including the $H$ dibaryon and other exotic hadronic systems with strangeness are studied in this context.

† talk presented at the YITP Workshop on “From Hadronic Matter to Quark Matter”, Kyoto, October 30 – November 1, 1994; to be published in Prog. Theor. Phys. Suppl.

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1. Introduction

Strangeness is the third lightest flavor whose mass is comparable to the QCD scale parameter, $\Lambda_{\text{QCD}}$. The chiral $SU(3)_L \times SU(3)_R$ seems powerful in studying the low energy hadron phenomena, although the large strange quark mass breaks the $SU(3)$ symmetry significantly. Because the QCD coupling strength depends on the energy scale, QCD is not quite flavor blind. In fact, the light quark dynamics is very different from that of heavy quarks. The heavy quark symmetry reveals that the light degrees of freedom can be separated from the heavy ones in many of heavy hadron phenomena.

Is the strange quark dynamics just a repetition of the light one with complication due to the symmetry breaking? I here argue that some new important and interesting physics issues arise from the introduction of the new flavor. I itemize them and will explain (most of) them later.

(a) The $U_A(1)$ anomaly represented by the instanton-induced interaction, which enhances the flavor mixing in the pseudoscalar mesons and plays significant roles in multi-hadron systems.

(b) The Pauli “allowance” in the multi-hadron systems or the matter, which makes the strange quark matter, the $H$ dibaryon, heavy multistrange nuclei plausible.

(c) The strangeness changing weak processes, which are experimentally studied in hypernuclear decays and reveal new features of hadronic weak interactions.

Strangeness plays major roles in various stages of the hadron physics.

(1) First, the single hadron spectrum is (historically) important in revealing the $SU(6)$ symmetry breaking and the chiral $SU(3)_L \times SU(3)_R$ symmetry. There the symmetry breaking is minimal as is seen for instance in the baryon magnetic moments. Namely, the $SU(3)$ or $SU(6)$ symmetry is good for the currents and wave functions for mesons and baryons. This makes the naive valence quark model successful. At the same time the explicit symmetry breaking and the $U_A(1)$ anomaly are found to play important roles in the single hadron system.

(2) When we turn to hadronic interactions, a new aspect of the third flavor becomes important. It is the Pauli exclusion principle. Strangeness in nonstrange environment is
much freer than the nonstrange quarks and thus the spectrum is richer than the nonstrange systems. One such example is the $H$ dibaryon, which is a strangeness $-2$ six-quark bound state. I will also show that the hyperon-nucleon interactions exhibit a variety of interesting features of hadronic interactions. Experimental efforts on the hyperon and hypernuclear physics have developed a new exciting field in hadron physics.

(3) In hadronic matter, various interesting possibilities have been suggested, such as, the strange quark matter, multistrange heavy nuclei, kaon condensation, K balls. Here again the symmetry breaking is important. Comparing to the pion, the kaon is relatively heavy. This makes some of the strange systems more stable than the corresponding nonstrange systems.

(4) Weak interactions of strangeness reveal new roles of hadronic interactions. The $\Delta I = \frac{1}{2}$ rule, for instance, is experimentally confirmed in various single hadron decays, and yet its mechanism is not fully understood. Interference between the standard theory of the weak interaction and the strong hadronic interaction is to be studied from this viewpoint.

In this article, I would like to review some of the above mentioned “strange” subjects in order.

2. Strangeness in Hadron

2.1. Instanton Induced Interaction

The instanton is a gluon field configuration with nontrivial topology and has provided a qualitative understanding of various nonperturbative features of QCD. One of its important consequences is that the coupling of light quarks ($u$, $d$ and $s$) to the instanton breaks the axial $U_A(1)$ symmetry and causes the $\eta$-$\eta'$ mass splitting.

Dynamical properties of instantons in the QCD vacuum have been studied by Shuryak and his collaborators. They found that the QCD vacuum can be described by an instanton liquid composed of small-size ($\simeq 0.3$ fm) instantons and anti-instantons. The density of instantons is low ($\simeq 1$ fm$^{-4}$) so that orientations of near-by instantons are almost independent from each other. This gives a gluon condensate consistent with the
QCD sum rule and a quark condensate as a result of dynamical breaking of the $SU(3)_L \times SU(3)_R$ chiral symmetry. The light quarks get constituent masses of a few hundred MeV due to the symmetry breaking.

A recent lattice QCD calculation tested the instanton liquid picture of the QCD ground state with the help of the cooling technique. After a cooling it has been shown that only a few instantons and anti-instantons survive and yet the meson and baryon correlators are essentially unchanged from the original uncooled calculation. It thus indicates that the instantons in the vacuum play the most essential role in determining the low lying mesons and baryons.

The $U_A(1)$ symmetry breaking is attributed to the instanton-light quark coupling. 't Hooft showed that a massless quark forms a zero mode around the instanton, which dominates the light-quark instanton coupling at low energy. In the instanton liquid vacuum, quarks form a zero mode around each instanton independently and propagate by hopping from one zero mode to another.

The effective interaction of light quarks in the instanton vacuum, called instanton induced interaction $\text{III}$, is given by a local effective interaction,

\[ H^{(3)} = V_0 \bar{q}_R(1)q_R(2) \bar{q}_R(3) \frac{189}{40} A_3^{(1)} \left(1 - \frac{1}{7} \sum_{i>j=1}^3 \vec{\sigma}_i \cdot \vec{\sigma}_j\right) q_L(3)q_L(2)q_L(1) + \text{(h.c.)} \]  

Here $q_{L(R)} \equiv \frac{1+\gamma_5}{2} q$ is the left (right) handed quark field operator, and $A_n^{(1)}$ is the projection operator to the flavor singlet state of $n$ quarks. The hermitian conjugate term (h.c.) arises from the interaction through the anti-instanton. The strength $V_0$ depends on the density of the instanton in the QCD vacuum.

The $\bar{R}$ (right-handed) – $L$ (left-handed) structure of eq.(1) clearly shows the chirality nonconservation. It however applies only to the flavor singlet part because $\text{III}$ is proportional to the projection $A_3^{(1)}$ into the flavor singlet states. Thus the $U_A(1)$ symmetry is broken by $\text{III}$, while the $SU(3)_L \times SU(3)_R$ is preserved.

The six-quark $\text{III}$ induces a four quark vertex $\text{III}_2$ through the contraction of a pair of $\bar{q}_R$ and $q_L$ by the current quark mass term or by the quark condensate in vacuum,

\[ H^{(2)} = V_0^{(2)} \bar{q}_R(1)\bar{q}_R(2) \frac{15}{8} A_2^{(1)} \left(1 - \frac{1}{5} \vec{\sigma}_1 \cdot \vec{\sigma}_2\right) q_L(2)q_L(1) + \text{(h.c.)} \]
The strength of the two-body interaction $V_0^{(2)}$ is related to the three-body strength and the quark condensate of the third flavor [5],

$$V_0^{(2)}(1, 2) \equiv \frac{1}{2} V_0(\langle \bar{q}q \rangle - K \tilde{m}_3) = \frac{1}{2} V_0(-Km_3)$$  (3)

where $K$ is a positive constant and $\tilde{m}_i (m_i)$ is the current (constituent) quark mass of the $i$-th flavor. The $SU(3)$ symmetry breaking is represented here by the strange quark mass so that the interaction between the $u - d$ quarks is stronger than that for $s - u$ or $s - d$.

The ratio, $V_0^{(2)}(u, s)/V_0^{(2)}(u, d)$, is given by $\xi \equiv m_u/m_s$. The hadron spectrum indicates $\xi \approx 0.6$, while the QCD sum rule suggests $\xi \approx 0.8$ [3]. Note that $III_2$ is attractive, while the full $III$, eq.(1), is repulsive, because $\langle \bar{q}q \rangle$ is negative.

2.2. Meson

Many studies have been carried out for the meson spectrum in the three-flavor Nambu-Jona-Lasinio (NJL) model [6]. This model is particularly useful in understanding the pseudoscalar mesons which are regarded as the Nambu-Goldstone (NG) bosons for the chiral symmetry breaking. There the 6-quark $III$, eq.(1), is commonly used to represent the $U_A(1)$ anomaly. For the $q\bar{q}$ meson system, the loop contraction of a $q\bar{q}$ pair leads to the 4-quark effective vertex $III_2$, eq.(3). As $III_2$ requires the projection to the flavor singlet (antisymmetric) state, it contributes to the $\eta_1$ meson repulsively, while $\pi$ and $\eta_8$ get strong attraction. Thus the $U_A(1)$ symmetry is broken explicitly by $III$.

The mixing of $\eta_1$ and $\eta_8$ comes mainly from the $SU(3)$ symmetry breaking. The large strange quark mass makes $(s\bar{s})$ separated from $(u\bar{u})$ and $(d\bar{d})$. Without the $U_A(1)$ breaking, the ideal mixing, i.e., $\eta' = s\bar{s}$ and $\eta = (u\bar{u} + d\bar{d})/\sqrt{2}$ would be realized, while $III_2$ mixes the $s\bar{s}$ with $(u\bar{u} + d\bar{d})/\sqrt{2}$. Thus the instanton induced interaction and the $SU(3)$ breaking due to the quark mass work in the opposite direction for the flavor mixing.

Due to the $R - L$ structure, the $III_2$, has the opposite sign for the scalar and pseudoscalar channels. It turns out that the large flavor mixing in the $\eta - \eta'$ system is consistent with the instanton induced interaction. Namely, the $III_2$ is strongly attractive in the pseudoscalar channel so that the chiral symmetry is dynamically broken and generates the $\langle \bar{q}q \rangle$ condensate, and the pseudoscalar octet mesons $\simeq (\bar{q}(\lambda_i/2)\gamma^5q)$ are transmuted
into the massless (light) NG bosons. In the scalar channel the flavor mixing is suppressed and therefore the OZI separation is achieved. The mixings in the vector and the axial vector channels are also suppressed to order \(1/N_c\). Thus, the OZI rule is well satisfied in the scalar, vector and axialvector channels, while the pseudoscalar mesons have a large flavor mixing. In conclusion, the light-quark-instanton interaction provides a consistent picture of the \(U_A(1)\) breaking and the flavor mixing in various channels.

### 2.3. Baryon

The baryon mass spectrum and the magnetic moments support the spin-flavor \(SU(6)\) symmetry. Also the meson-baryon couplings are consistent with the \(SU(3)\) (and \(SU(6)\)) symmetry of the wave functions. The explicit breaking of \(SU(3)\) is minimal and is taken care by the Gell-Mann-Okubo formalism, which treats the \(SU(3)\) octet term in the hamiltonian perturbatively.

Contrary to the mesons, which are deeply bound and thus are to be treated relativistically, the conventional constituent quark model gives a simple and good picture for baryons. Three valence constituent quarks confined in a potential well interact with each other by exchanging gluons. The one gluon exchange interaction contains a color-magnetic term, that breaks the \(SU(6)\) symmetry. A relevant term of OGE in the following discussion is the central part of the color magnetic interaction (CMI), which reads in the Breit-Fermi form

\[
V_{\text{CMI}} = -\frac{\alpha_s}{4} \sum_{i<j} \frac{2\pi}{3m_i m_j} (\lambda_i \cdot \lambda_j) (\vec{\sigma}_i \cdot \vec{\sigma}_j) \delta(\vec{r}_{ij}) \tag{4}
\]

The \(N - \Delta\) and the \(\Lambda - \Sigma\) mass difference both arise from this term, providing that the ratio of the strange and nonstrange quark masses are about 0.6. The same interaction deforms the nucleon wave function, which results in the nonvanishing neutron charge radius.

What are the roles of the instanton induced interaction in the baryon? For the valence quark model, we reduce III, eqs. (1) and (2), into the following nonrelativistic forms:

\[
V_{\text{III}} = V_0 \frac{189}{40} \sum_{(ijk)} A_j \left[ 1 - \frac{1}{7}(\sigma_i \cdot \sigma_j + \sigma_i \cdot \sigma_k + \sigma_j \cdot \sigma_k) \right] \delta(\vec{r}_{ij}) \delta(\vec{r}_{jk}) \tag{5}
\]
\[ V_{\text{III}2} = U_0^{(2)} \sum_{i<j} \frac{1}{m_im_j} \left[ 1 + \frac{3}{32} \lambda_i \cdot \lambda_j + \frac{9}{32} \bar{\sigma}_i \cdot \bar{\sigma}_j \lambda_i \cdot \lambda_j \right] \delta(\vec{r}_{ij}) \]  

(6)

where \( U_0^{(2)} = -\frac{1}{2} V_0 K m_0^2 m_s \), and \( m_i \) is the constituent quark mass.

It is interesting to note that the two-body instanton-induced interaction contains the identical term to eq.(4), and thus is not distinguishable from the one-gluon exchange as far as the hyperfine splittings are concerned. We recently proposed a view that both the instanton induced interaction and the one-gluon exchange interaction share the hyperfine splittings of a meson-baryon spectrum[6,7,8]. For simplicity, we fix the total strength of the spin-spin interaction so as to reproduce the nucleon delta mass difference, which is the typical hyperfine splitting in the baryon spectrum. We represent the share of the instanton induced interaction by \( p_{\text{III}} \), i.e.,

\[ H_{\text{hyp}} = p_{\text{III}} V_{\text{III}2} + (1 - p_{\text{III}}) V_{CMI}. \]

(7)

In order to determine \( p_{\text{III}} \), the \( \eta - \eta' \) mass difference is calculated. We find that \( p_{\text{III}} \approx 30 - 40\% \) gives a reasonable mass difference[8].

It is also important to notice that the three-body instanton-induced interaction has no direct contribution to the valence quarks in the baryon, if the small size instanton is assumed. This is the result of the symmetry consideration. The three valence quarks are color antisymmetric and must be flavor antisymmetric, orbital symmetric for the instanton induced interaction. But this combination is not allowed for the fermion.

This leads to the conclusion that contrary to the mesonic case the instanton induced interaction does not disturb the single baryon spectrum successfully explained in the conventional quark models.

3. Strangeness in \( B = 2 \)

When strangeness is planted into hadronic matter, a new important effect comes into effect. It is the Pauli principle. Witten first suggested that the strange quark matter (strangelet) can be the most stable system with finite baryon number, because the Pauli principle prefers mixing of strangeness in the ordinary matter[10]. The ratio of the
strangeness and the baryon number is of order one in the strangelet. Despite the heavier strange-quark mass, large strangeness mixing is favored because of the high Fermi energy for the ordinary quarks in bulk matter.

The quark antisymmetrization plays a major role even in $B = 2$ systems with strangeness. In fact, the reason why the deuteron binding energy is much smaller than the energy scale of ordinary hadronic interaction can be understood by the symmetry consideration. It is the strong short-range repulsion that prevents the deuteron, or any other nuclei, from crashing into quark matter. The simple quark model symmetry predicts that the Pauli principle with the help of the hyperfine interaction among quarks yields strong short-range repulsion in most of two-baryon systems\[11\]. Introduction of new flavor, the strangeness, may reduce the Pauli exclusion requirement, i.e., the Pauli allowance and thus makes 6-quark type dibaryons possible to exist. The most plausible candidate is the $H$ dibaryon.

3.1. $H$ DIBARYON

The $H$ dibaryon is a bound state of six quarks, or two baryons with strangeness $-2$. The strong hyperfine interaction due to the gluon exchange favors the flavor singlet 6-quark state\[12\], and indeed several quark model calculations predict a bound state, $H$, with $B = 2$, $J^{π} = 0^+$, and strangeness $-2$\[13,14\]. $H$ couples to the ordinary two baryon channels, $\Lambda\Lambda$, $N\Xi$, and $\Sigma\Sigma$. If the $H$ mass is below the $\Lambda\Lambda$ threshold, $H$ decays (into $N\Sigma$) only by the weak interaction.

Experimental searches of $H$ are underway in two facilities, BNL and KEK. The KEK group has found two candidates for double hypernuclei, $\Lambda\Lambda A$, which are formed in the $(K^+, K^-)$ experiment on the emulsion target. The estimated binding energy of the double hypernuclei places the upper limit of the $H$ binding energy below the $\Lambda\Lambda$ threshold. Thus, the possibility of deeply bound $H$ with binding energy 30 MeV or more has been excluded. Existence of $H$ should have a significant effect in the formation of double hypernuclei and $\Xi$-hypernuclei. Sakai et al. studied the interaction between $H$ and nuclei in the quark model and suggested possibility of $H$ bound weakly to nuclear matter\[15\].
When the binding energy of $H$ is small, we cannot neglect the effects of two baryon continuum states. The quark cluster model (QCM) calculation is more appropriate for this situation than the bag model, as one can couple the $\Lambda\Lambda$, $N\Xi$, and $\Sigma\Sigma$ channels with $H$. The QCM calculation for $H$ has been carried out by several groups\cite{13,14}. The results, which seem sensitive to the choice of the long-range meson exchange interaction, indicate that $H$ is stable with the binding energy ranging from 40 to 120 MeV unless the instanton effect is taken into account. We, however, found that $H$ as the six-quark system with strangeness has significant contribution from the 3-body instanton induced interaction and that $H$ in this channel is strongly repulsive\cite{7}. The effect is more than 50 MeV, and results in an unstable or barely bound $H$ dibaryon. It has also been found that the flavor non-singlet component $H'$ (with the flavor [42] symmetry) becomes almost degenerate with the flavor singlet $H$ for a strong instanton induced interaction\cite{10}.

We have to wait for further experimental data to make a conclusion, but it is plausible now that $H$ is not a deeply bound 6-quark state, but a weakly bound state of $\Lambda\Lambda$ and $N\Xi$. In fact, the flavor singlet combination is dominated by $N\Xi$ and the quark cluster model analysis supports the large $N\Xi$ contribution in $H$.

3.2. Hyperon-Nucleon Interaction

Study of hypernuclei has revealed that the hyperon bound in nuclear matter has a binding energy less than that of the nucleon. This fact clearly suggests that the hyperon-nucleon (YN) interaction is as strongly repulsive as the nuclear force at short distances, because the long-range part of the interaction should not be much weaker for YN. Low energy YN scattering data support the existence of the short range repulsion between YN.

The phenomenological YN potentials\cite{17,18} based on the $SU(3)$-symmetric meson exchange interaction usually assume a repulsive core or a form factor cutoff to represent the repulsion. The core radii or the cutoff masses are free parameters to be determined phenomenologically in each channel. Because the number of experimental data is so small that the fit is not always unique. Different potential models have sometimes very different
core radii.

On the other hand, the quark model description of the short-range repulsion between two nucleons is quite successful[19,20,11]. It has been demonstrated in the quark cluster model (QCM) calculations that the color-magnetic gluon exchange and the quark antisymmetrization in the valence quark dynamics provides a non-local soft repulsive core, which can reproduce the $N - N$ scattering S matrices for energies up to 300–400 MeV. When QCM is applied to other two-baryon systems, the same mechanism yields strong short-range repulsions in most of the two octet-baryon systems, including $\Lambda$-$N$ and $\Sigma$-$N$[21,11]. We find that the quark exchange effects (due to the antisymmetrization) show distinctive spin-isospin dependences especially for the $\Sigma$-$N$ interactions. Such strong channel dependences are attributed for the $SU(6)$ quark symmetry of the two-baryon system and have not been considered or taken into account in the conventional YN potential models. Thus it is quite interesting to see whether a realistic YN interaction that incorporates the quark-exchange mechanism can reproduce the YN two-body data as well as the hypernuclear spectrum.

Along this line of thought we construct a QCM-based potential model for the hyperon-nucleon interaction incorporating both the meson exchanges and the quark-gluon effects[22]. Such a model enables us to analyze experimental YN scattering data and to determine whether the quark-exchange mechanism is indeed at work for the hyperon-nucleon systems. This is important further for the study of double strange systems, such as $\Lambda$-$\Lambda$, $N$-$\Xi$ and the H dibaryon, because the interactions of $S = -2$ two-baryon systems are not yet directly accessible in experiment and thus require theoretical predictions.

3.3. Realistic YN Potential Model

We consider a valence quark model with a hamiltonian,

$$H = K + V_{\text{CONF}} + V_{\text{OGE}}$$  \hspace{1cm} (8)

*This section is based on the research carried out by K. Ogawa, S. Takeuchi and M. Oka.[22]*
where $K$ is the nonrelativistic quark kinetic energy term, $V_{\text{CONF}}$ stands for a quark confinement potential and $V_{\text{OGE}}$ is the Fermi-Breit potential for the one gluon exchange. We employ the resonating group method (RGM) wave function for the six-quark system, given by

$$\Phi_{BB'}(1 \sim 6) = \mathcal{A}[\phi_B(1 \sim 3) \phi_{B'}(4 \sim 6) \chi(R)] \quad (9)$$

and solve the RGM integral equation, with the kernels $H$ (Hamiltonian) and $N$ (Normalization):

$$\int [H(R, R') - E N(R, R')] \chi(R') dR' = 0 \quad (10)$$

Nonlocality of the RGM equation comes from the antisymmetrization of the quarks[19].

In order to describe the long-range part of the baryon-baryon interaction, we additionally need the meson exchange potentials. We keep the SU(3) symmetry for the meson-baryon couplings. Indeed, the $YN$ potential models, such as the Nijmegen models[17] and Jülich models[18], are based on the SU(3) symmetry. In this study, we employ a one-boson exchange potential based on the Nijmegen potential model and instead of using the hard cores in the original model, superpose it with the quark exchange interaction at the short distance. We introduce to the QCM equation (11) the meson exchange potential, which is of the form of the Nijmegen model. This can be done by adding an integral kernel for the meson exchange potential[20], given by

$$V(R, R') \equiv \int dR'' N^{1/2}(R, R'') V_f(R'') N^{1/2}(R'', R') \quad (11)$$

where $V_f$ is the meson exchange potential with the appropriate form factor. The form factor is chosen so as to be consistent with the quark wave function of the baryon,

$$V_f(R) \equiv \int \rho(x; R/2)V_{\text{meson}}(x - y)\rho(y; -R/2) \, dx \, dy \quad (12)$$

where $V_{\text{meson}}$ contains the exchange potentials of the pseudoscalar nonet, the vector nonet and the scalar mesons, whose coupling constants satisfy the $SU(3)$ relations. The form factor $\rho(x; R/2)$ is taken as the quark density of the baryon centered at $R/2$, normalized
\[ \int \rho(x) \, dx = 1 \]  

(13)

In the QCM calculation, we employ the Gaussian for the internal quark wave functions of the baryon for simplicity, and thus the corresponding form factor is given also by a Gaussian.

Our main strategy is (1) to fix the model parameters, such as the quark model parameters and the meson-baryon coupling constants, so as to reproduce the \( N - N \) scattering data and (2) generate the \( YN \) interaction using the \( SU(3) \) symmetry and its breaking due to the quark mass difference. We then (3) compare the \( YN \) scattering S matrices in our calculation with those from the conventional potential models, such as the Nijmegen models.

We find that the Pauli exclusion principle gives a stronger repulsion for the \( N\Sigma \) \((S = 0, I = \frac{1}{2})\) and \( N\Sigma \) \((S = 1, I = \frac{3}{2})\) channels, while the other channels show a mild repulsion which is generally softer than the original Nijmegen model D. The repulsion in the \( N\Sigma \) \((I = \frac{3}{2}, \frac{3}{2}S_{1})\) channel is as strong as that in the Nijmegen model F, which is known to provide not enough binding for \( \Sigma \) to make a bound \( \Sigma \) hypernuclei. The strong spin-isospin dependence of the \( N - \Sigma \) interaction is expected from the symmetry consideration in the quark cluster model[11]. The stronger repulsions for \((S = 0, I = \frac{1}{2})\) and \((S = 1, I = \frac{3}{2})\) are attributed to the Pauli exclusion of the six-quark state made of the completely overlapping \( N - \Sigma \) with these quantum numbers.

We also calculate the \( YN \) cross sections and the polarizations in order to study effects of the quark exchange repulsion to the observables. The low energy \( \Lambda - N \), and \( \Sigma - N \) scatterings have not been fully studied experimentally. A few available cross section data are compared with our results. They suggest that the scalar meson exchange, which gives the major part of the attraction, should contain both the flavor singlet and octet components and their mixing. Indeed, if we assume that the scalar \( \sigma \) meson with the mass about 550 MeV is flavor singlet, then the \( \Lambda - N \) and \( \Sigma - N \) cross sections at low energy are too much enhanced. Recent study of the \( \Sigma \) potential in our model by Takeuchi, Takayanagi and Shimizu indicate that the QCM potential with the flavor singlet \( \sigma \) meson
makes $\frac{5}{2}He$ bound too much [23]. Their conclusion is consistent with the hypernuclear spectroscopy, which indicates that the mean field potential for $\Lambda$ or $\Sigma$ is weaker than that for the nucleon.

Similar study of the $YN$ interaction has been carried out by the Tübingen group [24] and the Niigata–Kyoto group [25]. The Tübingen group proposed a QCM approach with the chiral scalar ($\sigma$) and the flavor-octet pseudoscalar mesons coupled directly to valence quarks. The $NN$ scattering phase shifts are fitted successfully only after the $\sigma NN$ coupling constants are modified for the $L \geq 1$ partial waves. They also introduce baryon mass dependences in the scalar-meson-exchange potential, which effectively reduces the medium-range attraction in the $YN$ interactions.

The Niigata–Kyoto group employs the quark cluster model with a scalar-nonet meson exchange and the tensor part of the $\pi$ and $K$ exchanges. They pointed out that the flavor $SU(3)$ symmetry is preserved in the global feature of the interaction and argue that the attraction due to the scalar meson exchange should be weaker for $\Lambda N$ and $\Sigma N$ so that the effect of larger reduced mass is compensated. Their potential model RGM–F based on the Nijmegen F potential can successfully reproduce the $NN$ phase shifts as well as the available $YN$ cross section data [25].

Dover and Feshbach studied the $SU(3)$ symmetry structure of the $YN$ interactions [26]. They found that both the long-range meson exchange potentials and the short-range quark-exchange repulsion contain the $SU(3)$ symmetry breaking primarily confined to the diagonal transitions. The available $YN$ cross sections seem to support their conclusion. This therefore indicates that the $SU(3)$ breakings in the meson and quark exchanges come naturally from the meson/quark mass differences.

Compared with the progress made in these theoretical programs, the experimental data for two-body $YN$ interactions are yet poor. While a recent measurement of the $\Sigma - N$ scattering cross section at KEK is under analysis [27], further experimental study of the two-body $YN$ systems is desirable.
3.4. Multistrange Matter

The third flavor increases the possibility of finding more exotic hadron systems. One of the examples is multistrange hypernuclei. Schaffner et al.\textsuperscript{[28]} predicted deeply bound multistrange nuclei with the strangeness fraction $|S|/A \simeq 1$ and the binding energy per baryon $E_B/A \geq 20$ MeV. Such systems contain many $\Xi$ baryons because the Pauli principle makes them stable against the strong decay, $\Xi N \rightarrow \Lambda \Lambda$.

Another strange hadron system of interest is K-ball\textsuperscript{[29,30]}, which is a nontopological soliton (Q-ball) in the chiral theory and is considered as a (classical) bound state of many kaons. The K ball with large strangeness may be stable against strong decay modes and thus be a long-lived mesonic object. Possibility of such object was first studied by Distler et al.\textsuperscript{[29]}, and indeed the solutions are later constructed in a specific model, the scale-invariant $SU(3)$ chiral effective theory\textsuperscript{[30]}. It was pointed out that the strangeness again favors such nontopological soliton solutions, mainly because the kaon is much heavier than the pion and therefore the kaon bound state may be more stable than pions. These exotic systems with multiple strangeness might be produced and detected in high energy heavy ion collisions. Indeed, it was suggested\textsuperscript{[31]} that coalescence of strange baryons in heavy ion collisions has a significant probability of producing the $H$ dibaryon if it is a bound state.

4. $\Delta S = 1$ Weak Transition

Free $\Lambda$ decays weakly into a nucleon and a pion with two isospin modes, $p\pi^-$ and $n\pi^0$, which share 64% and 36% of the total decay rate. If the decay goes through a $\Delta I = \frac{1}{2}$ vertex, the $p\pi^-/n\pi^0$ ratio would be two to one except for a small correction due to the phase space difference. The experimental ratio is very close to the $\Delta I = \frac{1}{2}$ prediction, and thus support the $\Delta I = \frac{1}{2}$ hypothesis for the hadronic weak decay.

In the nuclear medium, the $\Lambda \rightarrow N\pi$ decay is suppressed by the Pauli blocking on the final nucleon state, whose momentum is less than 100 MeV/c for the $\Lambda$ decay at rest. Indeed, in heavy hypernuclei, the decay is predominantly the nonmesonic one, that is, $\Lambda N \rightarrow NN$. If we assume that the initial $\Lambda$ and $N$ are at rest, then the final relative
momentum of $NN$ is about 420 MeV/c and thus is well above the Fermi momentum. Furthermore, this final momentum is high enough to look into the short-distance component of the two nucleon system and therefore this type of the hyperon decay may reveal a new aspect of the weak interaction of quarks under the influence of the strong interaction.

We here study the direct quark (DQ) processes in the two-body $\Lambda N \rightarrow NN$ weak decays and point out qualitative differences from the conventional picture, i.e., the meson ($\pi$, $K$, $\rho$, etc.) exchange mechanism, where one of the meson-baryon vertices involves the weak transition $s \rightarrow d$. It is quite timely to study this process in detail because new accumulating experimental data have revealed some difficulties in the meson-exchange picture.

For instance, the so-called $n-p$ ratio, i.e., the ratio $R_{np}$ of $\Lambda n \rightarrow nn$ v.s. $\Lambda p \rightarrow np$ decay in the nucleus, is predicted very small, $R_{np} \simeq 0.1 - 0.4$ in the meson-exchange picture. This is due to the strong contribution of the tensor force, which is preferred at the large momentum transfer. The tensor force selects the $S = 1$, $I = 0$ pn final state and therefore $R_{np}$ becomes small. The experimental data seem not to agree with the prediction, i.e., $R_{np}^{exp} \simeq 1$ in decays of light hypernuclei. We argue that the direct quark process, which does not follow the $I = 0$ selection rule, may enhance the $n-p$ ratio.

Another important aspect of the nonmesonic weak decays is the $\Delta I = \frac{1}{2}$ rule, which is known to be satisfied in the mesonic weak decays of hyperons to about 5% error. The same rule for the nonmesonic weak processes, like $\Lambda N \rightarrow NN$, is not confirmed yet.

Recently, Schumacher proposed to check the validity of the $\Delta I = \frac{1}{2}$ rule in the nonmesonic decays of the s-shell hypernuclei. It is therefore important to clarify the mechanism of the $\Delta I = \frac{1}{2}$ rule in the free hyperon decays and to study whether the same mechanism restricts the nonmesonic decays to $\Delta I = \frac{1}{2}$ as well.

4.1. Effective Weak Hamiltonian

The effective weak hamiltonian describing $\Delta S = 1$ processes has been calculated by several authors. It can be computed by evaluating perturbative QCD corrections using the operator product expansion and the renormalization group equation for the Wilson coefficients. It should be noted that new four-quark operators mix to the pure
weak \( su \to du \) vertex and thus change the flavor (isospin) structure significantly. The effective weak hamiltonian at the renormalization scale \( \mu^2 \) \((\leq 1 \text{ GeV}^2)\) is given by

\[
H_{\Delta I=1}^{\text{eff}} = -\frac{G_f}{\sqrt{2}} \sum_{r=1,r\neq4}^6 K_r O_r
\]

where the four-quark operators, \( O_k \) \((k=1, 2, 3, 5 \text{ and } 6)\) are defined by

\[
O_1 = (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta)_{V-A} - (\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A}
\]

\[
O_2 = (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta)_{V-A} + (\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A}
+ 2(\bar{d}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta d_\beta)_{V-A} + 2(\bar{d}_\alpha s_\alpha)_{V-A} (\bar{s}_\beta s_\beta)_{V-A}
\]

\[
O_3 = O_3(\Delta I = \frac{1}{2}) + O_3(\Delta I = \frac{3}{2})
\]

\[
O_3(\Delta I = \frac{1}{2}) = \frac{1}{3} \left[ (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta)_{V-A} + (\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A}
+ 2(\bar{d}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta d_\beta)_{V-A} - 3(\bar{d}_\alpha s_\alpha)_{V-A} (\bar{s}_\beta s_\beta)_{V-A} \right]
\]

\[
O_3(\Delta I = \frac{3}{2}) = \frac{5}{3} \left[ (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta)_{V-A} + (\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A}
- (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta d_\beta)_{V-A} \right]
\]

\[
O_5 = (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta + \bar{d}_\beta d_\beta + \bar{s}_\beta s_\beta)_{V+A}
\]

\[
O_6 = (\bar{d}_\alpha s_\beta)_{V-A} (\bar{u}_\beta u_\alpha + \bar{d}_\beta d_\alpha + \bar{s}_\beta s_\alpha)_{V+A}
\]

with

\[(\bar{u}_\alpha s_\alpha)_{V-A} \equiv (\bar{u}_\alpha \gamma^\mu (1 - \gamma_5) s_\alpha) \quad \text{etc.} \]

\( \alpha \) and \( \beta \) denote the color of quarks and the color sum is always assumed.

Among these operators, \( O_3 \) contains a part that induces the \( \Delta I = \frac{3}{2} \) transition, \( O_3(\Delta I = \frac{3}{2}) \), while the others are purely \( \Delta I = \frac{1}{2} \). The coefficients \( K_r \) can be calculated by solving the renormalization group equation to the one-loop QCD corrections. It is important to note that the QCD correction enhances the \( O_1 \) coefficient and thus the \( \Delta I = \frac{1}{2} \) component, while the purely weak four-quark vertex \( su \to du \) contains both the \( \Delta I = \frac{1}{2} \) and \( \frac{3}{2} \) components. Later we will compare the results with and without \( O_3(\Delta I = \frac{3}{2}) \) in order to study the \( \Delta I = \frac{3}{2} \) contribution. In the present calculation, we
Table 1: Strengths of the weak effective four-fermi vertices, taken from ref. [36]. We use the version with flavor-number dependent Λ and \( m_t = 200 \text{GeV} \). The values of the CKM matrix elements are taken as the central values of those given in ref. [37].

| \( \mu \) (GeV) | \( \Lambda^{(4)} \) (GeV) | \( K_1 \) | \( K_2 \) | \( K_3 \) | \( K_5 \) | \( K_6 \) |
|----------------|----------------|--------|--------|--------|--------|--------|
| 0.24           | 0.10           | −0.284 | 0.009  | 0.026  | 0.004  | −0.021 |

employ a set of parameters from ref. [36] with \( \mu = 0.24 \text{ GeV} \), which is determined so as to give \( \alpha_s(\mu^2) = 1 \) for \( \Lambda_{QCD} = 0.1 \text{ GeV} \). The values of the coefficients \( K_r \) used in the present calculation are given in Table 1.

This effective Hamiltonian has been used for the calculations of the nonleptonic decay of strange mesons and baryons [33]. It is found that although the \( \Delta I = \frac{1}{2} \) transition is indeed enhanced in those decays, the enhancement is not enough to account for the experimental data quantitatively. It was suggested [38] that an additional \( \Delta I = \frac{1}{2} \) enhancement arises from the mesonic correction in the chiral effective theory. We have confirmed their conclusion by using the Nambu-Jona-Lasinio model without bosonization [39].

4.2. The \( \Lambda N \rightarrow NN \) Transitions

The effective weak Hamiltonian for quarks can be applied to the valence quark model of two-baryon systems or six-quark systems. We consider the direct quark (DQ) processes, i.e., the short-distance weak interaction of two-baryons without mediated by mesons. In calculating the DQ amplitudes for \( \Lambda N \rightarrow NN \), we employ the constituent quark model, which describes the spin-flavor structure of the ground-state baryons fairly well. Two-baryon systems are expressed by the quark-cluster-model wave functions [19,21]:

\[
\begin{align*}
|\Lambda N\rangle &= \mathcal{A}^6|\phi(1,2,3)\phi(4,5,6)\chi_i(\vec{R})\rangle \\
|NN\rangle &= \mathcal{A}^6|\phi(1,2,3)\phi(4,5,6)\chi_f(\vec{R})\rangle 
\end{align*}
\]

†This section is based on the research carried out by T. Inoue, S. Takeuchi and M. Oka [40].
where $\mathcal{A}^6$ is the antisymmetrization operator for six quarks, $\phi$ is the internal wave function of the baryon, and $\vec{R}$ is the relative coordinate of two baryons. $\chi_i(\vec{R})$ ($\chi_f(\vec{R})$) is the initial (final) relative wave function. We assume that the orbital part of $\phi$ is a Gaussian with the extension parameter $b = 0.5$ fm, and the flavor-spin part of $\phi$ is purely $\text{SU}(6)$ symmetric.

We here concentrate on the $s$-shell hypernuclei and assume that the initial state is represented purely by $L = 0$ and its wave function is given by

$$\chi_i(\vec{R}) = N_0 g(\vec{R}) \exp \left\{ -\frac{1}{2B^2} \vec{R}^2 \right\} \quad (22)$$

The final wave function is taken as

$$\chi_f(\vec{R}, \vec{K}) = g(\vec{R}) \exp \left\{ i \vec{K} \cdot \vec{R} \right\} \quad (23)$$

where $g$ represents the short range correlation,

$$g(\vec{R}) = 1 - C \exp \left[ -\frac{\vec{R}^2}{r_0^2} \right] \quad (24)$$

We choose the parameter $B$ as $\sqrt{2} \times 1.3$ fm for the $S$-shell hypernuclei, which corresponds to the shell model wave function with the Gaussian parameter $1.3$ fm for both the nucleon and $\Lambda$. For the short range correlation we choose $C = 0.5$ and $r_0 = 0.5$ fm in the present calculation. The relative momentum $K$ of the final state is determined by the realistic $Q$ value $\Delta E \equiv M_\Lambda - M_N$ of the decay: $K = 416$ MeV/c.

The results of the calculation are summarized in Table 2. As we restrict the initial state to $L = 0$, nine amplitudes, $a - f$ give all the information for the $\Lambda N \rightarrow NN$ weak decay. Amplitudes, $a$, $b$, and $f$ have contributions both from $\Delta I = 1/2$ and $\Delta I = 3/2$ transitions. The ratios $a_n/a_p$, $b_n/b_p$, and $f_n/f_p$ for $\Delta I = 1/2$ are equal to $\sqrt{2}$. The other amplitudes, $c$, $d$ and $e$ do not contain any $\Delta I = 3/2$ component, because the final $NN$ states have $I = 0$. We find that the amplitudes $a$ and $b$ get large $\Delta I = 3/2$ contributions, while $f$ is dominated by $\Delta I = 1/2$. Thus we conclude that the $\Delta I = 3/2$ transition can be studied in the weak decay starting from the $\Lambda N \ ^1S_0$ state.

The two-body transition rates are calculated for the $\Lambda N \rightarrow nN$ ($N = n$ or $p$) transitions with angular momentum $J$, and are denoted by $\Gamma_{NJ}^{\pi \hbar}$. Table 3 gives the results
Table 2: Calculated transition amplitudes in $10^{-10} \, MeV^{-1/2}$. This table is taken from ref. [40].

| Channel          | Direct Quark (DQ) | OPE  |
|------------------|-------------------|------|
|                  | full $\Delta I = 1/2$ | $\Delta I = 3/2$ |      |
| $p\Lambda \to pn$ |                   |      |      |
| $a_p \quad ^1S_0 \to ^1S_0$ | $-78.1$ | $-23.4$ | $-54.7$ | $2.2$ |
| $b_p \quad \to ^3P_0$   | $-53.5$ | $2.0$ | $-55.5$ | $-24.8$ |
| $c_p \quad ^3S_1 \to ^3S_1$ | $-1.0$ | $-1.0$ | $0$ | $2.2$ |
| $d_p \quad \to ^3D_1$   | $0$ | $0$ | $0$ | $-86.8$ |
| $e_p \quad \to ^1P_1$   | $-23.2$ | $-23.2$ | $0$ | $-43.0$ |
| $f_p \quad \to ^3P_1$   | $-55.4$ | $-53.8$ | $-1.5$ | $20.2$ |
| $n\Lambda \to nn$     |                   |      |      |
| $a_n \quad ^1S_0 \to ^1S_0$ | $5.5$ | $-33.1$ | $38.6$ | $3.1$ |
| $b_n \quad \to ^3P_0$   | $42.2$ | $2.9$ | $39.3$ | $-35.1$ |
| $f_n \quad ^3S_1 \to ^3P_1$ | $-75.1$ | $-76.2$ | $1.0$ | $28.6$ |

for $\Gamma_{NJ}$ as well as the spin averaged transition rates, defined by

$$\Gamma_{\Lambda p \to pn} = \frac{1}{4} (\Gamma_{p0} + 3\Gamma_{p1})$$  \hspace{1cm} (25)

$$\Gamma_{\Lambda n \to nn} = \frac{1}{4} (\Gamma_{n0} + 3\Gamma_{n1})$$  \hspace{1cm} (26)

We are also interested in the n–p ratio defined by

$$R_{np} \equiv \frac{\Gamma_{\text{neutron induced}}}{\Gamma_{\text{proton induced}}}$$  \hspace{1cm} (27)

We find that the $J = 0$ proton-induced transition rate, $\Gamma_{p0}$, is strongly enhanced due to the $\Delta I = 3/2$ transition. The results of the one-pion exchange (OPE) transition are also shown in Tables 2 and 3, for comparison [32]. These amplitudes satisfy $a_n/a_p = b_n/b_p = f_n/f_p = \sqrt{2}$, because $\Delta I = \frac{1}{2}$ is assumed for the weak pion-baryon vertex. One sees that the amplitude $d$ in this case is dominant. This comes from the tensor part of the one pion
Table 3: Decay rates of light hypernuclei. All the decay rates are in the unit of $\Gamma_{free}$. The experimental data for $\Gamma_{\Lambda p \rightarrow pn}$, $\Gamma_{\Lambda n \rightarrow nn}$, $\Gamma_{nm}(\Lambda\text{He})$ and $R_{np}(\Lambda\text{He})$ are taken from ref.[42]. Those for $R_{np}(\Lambda\text{He})$, $\Gamma_{nm}(\Lambda\text{He})/\Gamma_{nm}(\Lambda\text{H})$ and $\Gamma_{n0}/\Gamma_{p0}$ are taken from ref.[43]. This table is taken from ref.[40].

|                  | Direct Quark | OPE | DQ ± OPE | Exp       |
|------------------|--------------|-----|----------|-----------|
|                  | Full $\Delta I = \frac{1}{2}$ |      | +        | −         |
| $\Gamma_{p0}$    | 0.177        | 0.010 | 0.012    | 0.235     | 0.143     |
| $\Gamma_{p1}$    | 0.071        | 0.067 | 0.193    | 0.260     | 0.269     |
| $\Gamma_{n0}$    | 0.035        | 0.021 | 0.024    | 0.002     | 0.118     |
| $\Gamma_{n1}$    | 0.111        | 0.114 | 0.016    | 0.042     | 0.212     |
| $\Gamma_{\Lambda p \rightarrow pn}$ | 0.097 | 0.053 | 0.143    | 0.253     | 0.238     | 0.105±0.035 |
| $\Gamma_{\Lambda n \rightarrow nn}$ | 0.092 | 0.091 | 0.018    | 0.032     | 0.189     | 0.100±0.055 |
| $\Gamma_{nm}(5\Lambda\text{He})$ | 0.378 | 0.295 | 0.333    | 0.573     | 0.854     | 0.41 ± 0.14 |
| $R_{np}(5\Lambda\text{He})$ | 0.94 | 1.70 | 0.12     | 0.12      | 0.79      | 0.93 ± 0.55 |
| $R_{np}(4\Lambda\text{He})$ | 0.18 | 0.20 | 0.08     | 0.004     | 0.24      | 0.18 ± 0.12 |
| $\Gamma_{nm}(4\Lambda\text{He})/\Gamma_{nm}(4\Lambda\text{H})$ | 0.63 | 0.66 | 6.58     | 1.69      | 1.14      | 1.65 ± 0.77 |
| $\Gamma_{n0}/\Gamma_{p0}$ | 0.20 | 2.00 | 2.00     | 0.01      | 0.854     | −0.8 ± 2.7 |
| $a_1(5\Lambda\text{He})$ | 0.01 | 0.02 | −0.19    | 0.20      | −0.44     |

exchange and is enhanced due to a large relative momentum in the final state. Because the amplitude $d$ is not allowed for the $n\Lambda \rightarrow nn$ by the Pauli principle, the $n - p$ ratio in the pion exchange amplitudes becomes very small. The DQ amplitude $d$ is neglected because the tensor part of DQ is of order $(p/m)^2$.

Compared with the OPE result, $\Gamma_{p0}$ for DQ is much larger and in fact is dominant while OPE is dominated by the tensor transition included in $\Gamma_{p1}$. This dominance of $\Gamma_{p1}$ in OPE makes the $n-p$ ratio, $R_{np}$, small. In DQ, $\Gamma_{n1}$ is also large so that the spin averaged $R_{np}$ is as large as 1. Thus we find that DQ and OPE predict qualitatively different values for $R_{np}(\Lambda\text{He})$, while they give similar nonmesonic decay rate, $\Gamma_{nm}(\Lambda\text{He})$. The experimental data prefers DQ, which indicates a significant contribution of $\Gamma_{n1}$. Because the pure
\[ \Delta I = \frac{1}{2} \] calculation also yields large \( R_{np} \), its enhancement is not related to the \( \Delta I = \frac{1}{2} \) rule violation.

Recently, Schumacher\[33\] suggested that non-mesonic decays of the \( A=4 \) and 5 hypernucliei may be useful in checking the validity of the \( \Delta I = \frac{1}{2} \) rule. One can parametrize the n-p ratios of the decay of \( ^4 \Lambda \)He and \( ^5 \Lambda \)He and the ratio of non-mesonic decay widths of \( ^4 \Lambda \)He and \( ^4 \Lambda \) in terms of \( \Gamma_{NJ} \)[41]. Then, using experimental data, the ratios of \( \Gamma_{NJ} \) can be extracted. The \( \Gamma_{n0}/\Gamma_{p0} \) ratio will be 2 for the pure \( \Delta I = \frac{1}{2} \) transition, while it becomes 1/2 for the pure \( \Delta I = \frac{3}{2} \) transition. Our full amplitudes indeed predict small \( \Gamma_{n0}/\Gamma_{p0} \), which indicates a large \( \Delta I = \frac{3}{2} \) contribution, while the present experimental data do not have enough precision to be conclusive. It could be the first clear evidence for the \( \Delta I = \frac{3}{2} \) weak transition that is expected in the standard theory.

So far we have not considered the interference of the DQ and OPE amplitudes. The present formalism allows us to regard OPE independent from DQ and therefore to superpose these two amplitudes. Because the relative phase of two amplitudes are not known, we evaluate DQ \( \pm \) OPE and the results are given in Table 3. One finds that the difference between the two choices of the relative phase mostly appear in the neutron-induced decay rates. \( \Gamma_{nJ} \)'s are suppressed in (DQ + OPE) and thus the \( n - p \) ratio \( R_{np} \) becomes very small. In this sense, the experimental data prefer the (DQ – OPE) combination. The ratio \( \Gamma_{n0}/\Gamma_{p0} \) tends to be small (\( \ll 2 \)) for both (DQ \( \pm \) OPE) and again indicates a large \( \Delta I = 3/2 \) contribution. In both (DQ \( \pm \) OPE), we find that \( \Gamma_{\Lambda p \rightarrow p n} \) is overestimated and therefore the total nonmesonic decay rate \( \Gamma_{nm}(^4 \Lambda \)He) is too large. This is again due to the large tensor component in \( \Gamma_{p1} \) (OPE). It seems important that OPE is calculated in the quark interaction point of view in order to make a reliable prediction for the DQ – OPE interference.

5. Conclusion

I have reviewed several hadronic phenomena involving strangeness. I am particularly interested in the roles of the instanton induced interaction in the hadron and multibaryon spectra, the YN interactions from the quark model point of view and the strangeness
changing weak interaction among quarks. These subjects are studied in simple quark models, which give intuitive pictures of the underlying physics and further give quantitative understanding of experimental data. The study of strange hadron systems is still far to go for complete understanding and further effort both in experiment and theory is anticipated. Present experimental facilities have already achieved a significant progress in hypernuclear physics and are expected more results. New facilities such as CEBAF and RHIC are also expected to shed new light on the strangeness physics.

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