Quarks, Gluons and Weak Bosons in Hypernuclei

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Roles of the quark structure of baryons and mesons in hypernuclear physics are reviewed. Quark cluster model is employed to describe short-range part of hyperon-nucleon and hyperon-hyperon interactions. Important contributions of instanton induced interactions are stressed. New decay modes of Λ, $\Lambda N \to NN$, in hypernuclei are studied. Direct quark mechanism is introduced and possibility of breaking the $\Delta I = 1/2$ rule is discussed

§1. Why are we interested in strangeness?

In 2003, we celebrated the 50th anniversary of the discovery of the hypernucleus. The physics of strange hadrons and nuclei has flourished during these 50 years. The year 2003 is also marked by discoveries of the pentaquark and other narrow (exotic) hadrons with heavy quarks. It is quite timely, therefore, to summarize achievements in strangeness hadron physics to this date and reveal unsolved problems. That is the main purpose of the 18th Nishinomiya-Yukawa Memorial Symposium on "Strangeness in Nuclear Matter".[1)]

Why are we interested in strangeness? First of all, one sees a very rich hadron and nuclear spectrum with strange quarks. Those can be experimentally accessed by the K, Λ, Σ, and Ξ - nuclear reactions, which produce $S = -1$ systems, Λ, Σ, and K nuclei, $S = -2$ systems, $\Lambda\Lambda$, and Ξ nuclei, and $S = +1$ systems, Θ^+ (pentaquark). I would like to point out also that those strange systems help us to understand low-energy dynamics of the quantum chromodynamics (QCD).

A few important roles of the strangeness in hadron physics are enumerated.

1. Strangeness as impurity: The strangeness may go deep inside the u/d nuclear matter and may probe properties of (strange) hadrons in dense nuclear medium. It may create new forms of nucleus or matter, such as compact nuclei with a bound kaon proposed by Akaishi, Yamazaki and Doté.[2)] Recently, Iwasaki et al.[3)] reported a discovery of a deep bound kaonic nucleus, ${}^{3}_{K}H$ at the binding energy of about 170 MeV.
2. m_s vs QCD scale parameter: The pure glue classical QCD is scale invariant, but a scale parameter $\Lambda_{\rm QCD} \sim 210$ MeV is acquired by the scale (trace) anomaly. Because $m_s \sim 90$ MeV is of the order of $\Lambda_{\rm QCD}$, the strangeness is most sensitive to the QCD dynamics. We therefore expect that phenomena under the strong interaction depend on m_s in nontrivial ways.
3. Strangeness appears in compact stars: The neutron stars, and newly reported quark stars, are hyper-heavy-nuclei with strangeness. Takatsuka and Kaplan explained in the Symposium that the strange hadrons and their interactions change the equation of states of the high density nuclear matter, where there may exist kaon condensation and mixtures of the Λ, Σ, and Ξ hyperons. It is therefore critically important to understand the KN, YN and YY interactions

precisely.

In this article, I first introduce the current understanding of structures of hadrons and hadronic interactions from the quark model viewpoint. I particularly stress important roles of the spin-spin and the spin-orbit interactions between quarks. The interactions between the hyperon (Y) and nucleon (N) are studied in the quark model.

The second part is devoted to roles of quark substructure of the baryon in the weak decays of strangeness. The weak decay has been recognized as a very useful tool to reveal the symmetries and dynamical properties of the underlying dynamics of the elementary particles. In the hadronic weak interaction, we still have a few unsolved problems and the hypernuclear weak decay may help to solve some of them. An example is a so-called "$\Delta I = 1/2$ rule", whose mechanism has not been fully understood. We discuss possibility of violation of the $\Delta I = 1/2$ rule in weak decays of hypernuclei.

At the Symposium, I reviewed theoretical status of the pentaquark baryons. In this article, I omit that part and refer the readers to another review article.[4)]

§2. Quark model

The quark model is a powerful tool to understand hadron spectrum, structures and dynamics. It is, however, noted that the quark in the QCD Lagrangian is the current quark, which is, in principle, different from the constituent quark employed in the quark model. The constituent quark is considered as an effective fermion field that is relevant for structures and reactions of low energy hadrons. One way of defining and studying the constituent quark is to employ the Dyson-Schwinger equation (Fig. 1) for the quark propagator in QCD. Effective quark propagator is obtained as a solution of the DS equation. In this process, we note that the conserved currents are not renormalized. Therefore the isospin, I, the hypercharge, Y, the color C of the quark do not change due to the renormalization. The constituent quarks acquire effective quark masses,

$$\begin{aligned} m_q &\sim 350 \text{ MeV for the } u\text{, and } d \text{ quarks,} \\ m_s &\sim 500 \text{ MeV for the } s \text{ quark.} \end{aligned} \tag{2·1}$$

It is assumed that the residual interactions are weak after the quark mass is renormalized.

The baryon structure can be described in terms of the configurations of three

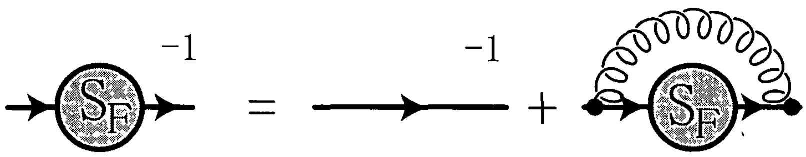

Fig. 1. A graphical representation of the Schwinger-Dyson equation for the quark propagator in the ladder approximation of QCD.

constituent quarks occupying single particle states. We naturally assume that the lowest-energy single particle state is an orbit with $L = 0$ and $S = 1/2$, labeled by $s_{1/2}$. The ground state baryons are then given by the configurations, $(s_{1/2})^3$ with ($J = 1/2$, $F = 8$) and ($J = 3/2$, $F = 10$). These correspond to the octet spin 1/2 baryons (N, Λ, Σ and Ξ), and the decuplet spin 3/2 baryons (Δ, Σ^*, Ξ^* and Ω).

The observed spectrum follows this classification, but one sees a significant mass splitting among them. The mass difference seems to have two origins; (1) the one responsible for the octet-decuplet splitting, and (2) the one for the splitting within the multiplet. The former is considered to come from the spin-spin interaction between the quarks, called the hyperfine interaction, while the latter is due to the SU(3) breaking caused by the heavy mass of the s quark.

2.1. *Hyperfine interaction*

The origin of the hyperfine interaction in the context of the color gauge theory of the strong interaction was first proposed by DeRujula, Georgi and Glashow.[5)] They considered exchange of one gluon between quarks and derived a spin-spin interaction arising from the magnetic part of the gluon (Fig. 2(a)). The spin-spin interaction is given in the nonrelativistic form by

$$\frac{\alpha_s}{m_i m_j}(\lambda_i \cdot \lambda_j)(\vec{\sigma}_i \cdot \vec{\sigma}_j)\delta(\vec{r}_{ij}), \tag{2·2}$$

which is called the color magnetic (CM) interaction.

For the ground state baryons (and the other states with s wave quarks), we may simplify the form of the CM by assuming that the orbital matrix elements are to the leading order the same, and obtain

$$\Sigma_{\rm CM} = \sum_{i<j} \Delta_{\rm CM}\, \xi_{ij}\, (\lambda_i \cdot \lambda_j)(\sigma_i \cdot \sigma_j). \tag{2·3}$$

We determine the strength $\Delta_{\rm CM}$ using the $N - \Delta$ mass difference:

$$\Delta_{\rm CM} = 18.75 \text{ MeV}. \tag{2·4}$$

In Eq. (2·3), ξ denotes a factor due to the SU(3) symmetry breaking, defined by

$$\xi_{ij} = \begin{cases} 1 & \text{for } (ij) = (uu),\ (ud) \text{ or } (dd), \\ m_u/m_s & \text{for } (ij) = (us) \text{ or } (ds), \\ (m_u/m_s)^2 & \text{for } (ij) = (ss). \end{cases} \tag{2·5}$$

This factor represents the fact that the $s-u$ and $s-d$ (and also $s-s$) CM interactions are weaker than that of $u - d$, and is important to explain the mass splitting (~ 80 MeV) of the Λ and Σ.

$$\Sigma_{\rm CM} = 150 \text{ MeV } \times \left(\frac{1}{3}(\vec{\sigma}_u \cdot \vec{\sigma}_d) + \frac{\xi}{3}(\vec{\sigma}_s \cdot (\vec{\sigma}_u + \vec{\sigma}_d))\right), \tag{2·6}$$

where we denote the mass ratio by $\xi \equiv m_u/m_s$. Then, for Λ with $(ud)I = 0, S = 0$, the expectation value of $\Sigma_{\rm CM}$ is 150 MeV $\times[(-1) + 0 \cdot \xi] = -150$ MeV, while for

Σ with $(ud)I = 1, S = 1$, 150 MeV $\times[(1/3) + (-4/3)\cdot\xi] = -80$ MeV, assuming $\xi \sim 3/5$. Thus the CM interaction causes the $\Lambda - \Sigma$ mass splitting of 70 MeV.

The CM interaction happens to be fairly successful in explaining the mass spectrum of the mesons and the baryons. It, however, has some shortcomings. One of the problems is that the coupling constant α_s has to be unreasonably large so that the perturbative expansion is not justified. For instance, in the MIT bag model, the value of α_s that is required in order to reproduce the baryon spectrum is as large as 2.

Another problem is the prediction of the H dibaryon. H is a six-quark state composed of $u^2d^2s^2$ quarks. In 1977, Jaffe[6] pointed out that the CM interaction causes a strong attraction in the flavor singlet 6 quark system and predicted a deeply bound H dibaryon.[7] The two-baryon threshold for this channel is $2M_\Lambda = 2230$ MeV. A crude estimate of the mass of H under the influence of the CM interaction is given by

$$\begin{aligned} M_H &= 6m_q + 2\Delta m + \langle \Sigma_{\rm CM}\rangle_H \\ &= 350\times 6 + 2\times 200 - 450 \sim 2100 \text{ MeV}, \end{aligned} \tag{2·7}$$

where we choose the u, d quark mass to be $m_q = 350$ MeV, and the mass difference of $\Delta m \equiv m_s - m_q = 200$ MeV, which are consistent with the masses of the N, Δ and Λ, i.e.,

$$\begin{aligned} M_N &= 3m_q + \langle \Sigma_{\rm CM}\rangle_N = 350\times 3 - 150 = 900 \text{ MeV}, \\ M_\Delta &= 3m_q + \langle \Sigma_{\rm CM}\rangle_\Delta = 350\times 3 + 150 = 1200 \text{ MeV}, \\ M_\Lambda &= 3m_q + \Delta m + \langle \Sigma_{\rm CM}\rangle_\Lambda = 350\times 2 + 550 - 150 = 1100 \text{ MeV}. \end{aligned} \tag{2·8}$$

However, searches of the H dibaryon for these 20 years were unsuccessful, and deny existence of a deep bound state. Later, Takeuchi and Oka[8] introduced a new interaction which happens to be repulsive in the H and thus explains non-existence of H. This is the subject of the next section.

2.2. *Instanton-induced-interaction (III)*

An alternative way of explaining the hyperfine interaction in QCD comes from a nonperturbative gluon configuration, called instanton. The instanton in the QCD vacuum induces a new type of interaction. (Fig. 2(b))

The instanton is a classical solution of the Yang-Mills field equation in the Euclidean 4-dimensional space. The solution has non-trivial topology and it interpolates QCD vacua with different topological numbers via quantum tunneling processes. The instanton-(light) quark coupling was first introduced by 't Hooft,[9] who pointed out that there exist zero modes of quarks localized around the instanton and that light quarks couple strongly to the instanton through those zero modes. He also pointed out that the instanton-induced-interaction (III), which is also known as the Kobayashi-Maskawa-'t Hooft (KMT) interaction, breaks unwanted axial $U_A(1)$ symmetry of the light quark sector of QCD. The $U_A(1)$ symmetry breaking manifests itself in the spectrum of pseudoscalar mesons as the mass splittings of $m'_\eta \gg m_\eta, m_\pi$. The III also causes mixings of $q\bar{q}$ with different flavors, and therefore is responsible

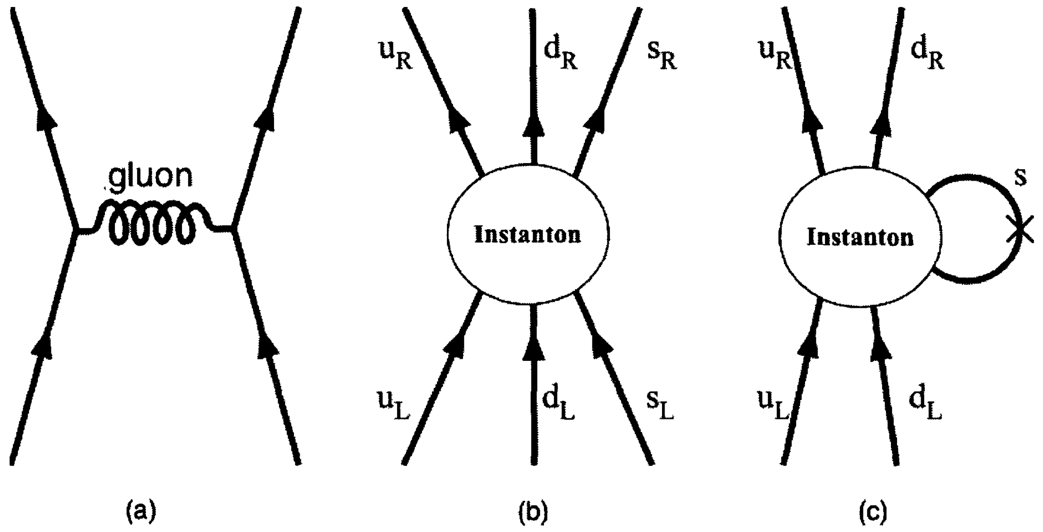

Fig. 2. Quark-quark interaction diagrams (a) for the one gluon exchange interaction, (b) for the three-body III and (c) for the two-body III, where the effective quark mass ($\times$) is inserted to make a contraction of the s quark.

for a large OZI breaking ($u\bar{u}$, $d\bar{d} \leftrightarrow s\bar{s}$) observed in the spectra of the pseudoscalar and scalar mesons.

The role of the instanton in the low energy hadron spectrum and reactions is further developed into the instanton vacuum picture,[10] where the path integral of the QCD vacuum is dominated by instantons. The instanton vacuum is consistent with the gluon condensate derived from the QCD sum rule, $\langle (G^{\mu\nu})^2 \rangle \sim 1$ instanton/fm^4. It was also pointed out that the III induces dynamical chiral symmetry breaking.

The III has a specific flavor structure. It is effective only on flavor antisymmetric combinations of quarks, such as, $u-d$ with $I=0$, and $u-d-s$ in the flavor singlet state. The range of the interaction is given by the size of the instanton, which is considered to be around 0.3 fm. Therefore the III is usually treated approximately as a contact interaction. The number of flavors participating in the III is given by the number of (nearly) zero modes around the instanton. The exact zero mode appears only when the quark is massless, but light quarks that contribute chiral symmetry breaking of QCD vacuum, should have nearly zero modes. For $N_f = 3$, the III is written as a three-body force among $u-d-s$, while the flavor SU(2) force is obtained from the three-body III by contracting a pair of quark and antiquark of the same flavor (Fig. 2(c)). The contraction is given either by the quark mass term or by the quark condensate in the QCD vacuum. The explicit forms of the three-flavor and two-flavor III are given by[11)]

$$V_{\rm III}^{(3)} = V^{(3)} \sum_{(ijk)} \mathcal{A}^f \left[1 - \frac{1}{7} \left(\vec{\sigma}_i \cdot \vec{\sigma}_j + \vec{\sigma}_j \cdot \vec{\sigma}_k + \vec{\sigma}_k \cdot \vec{\sigma}_i \right) \right] \delta(\vec{r}_{ij}) \delta(\vec{r}_{jk}),$$

$$V_{\rm III}^{(2)} = V^{(2)} \sum_{i<j} \mathcal{A}^f \left[1 - \frac{1}{5} (\vec{\sigma}_i \cdot \vec{\sigma}_j) \right] \delta(\vec{r}_{ij}). \qquad (2\cdot 9)$$

In the baryon spectrum, it is important that the two-flavor III contains a spin-

spin interaction. In 1989, Shuryak-Rosner[12] and Oka-Takeuchi[11] pointed out that the III has the right properties to explain the baryon spectrum: (1) The 2-body spin-spin interaction is attractive for flavor antisymmetric states, while it is zero for flavor symmetric states. Thus it is responsible for the N (octet)$-\Delta$(decuplet) mass splitting. (2) The strength of the spin-spin interaction depends on the constituent quark masses:

$$V_{su} \simeq \frac{m_u}{m_s} V_{ud}. \qquad (2\cdot10)$$

This mass dependence is exactly the one necessary for the baryon spectrum.

§3. Short range BB interaction

It is expected and also anticipated that hadronic interactions can also be explained from the constituent quark model points of view. In particular, it must be straightforward to apply the quark model to baryon-baryon interactions, because no antiquark is involved and therefore no annihilation diagrams exist. Quark model description of the nuclear force has been explored by many groups. The first idea came from Neudatchin, Smirnov and Tamagaki in 1977.[13] They proposed to explain the strong short-range repulsion between two nucleons by the Pauli exclusion principle of the quarks. It was later pointed out by Oka and Yazaki,[14] and also by Shimizu et al.[15] that the Pauli principle alone is not enough, but the color-magnetic interaction is necessary to obtain sufficient repulsion. The quark exchange diagram (Fig. 3(a)) gives strong repulsion between two nucleons. The range of the interaction is short because it requires two nucleons to overlap each other. The strength of the short-range repulsion is found to be of the order of the $N-\Delta$ mass difference $\simeq 300$ MeV (times an order-one constant).

The quark exchange interaction, however, does not yield the NN medium-range attraction that is responsible for the nuclear binding energy. In order to explain the medium and long range parts of the NN interaction, it was proposed to superpose conventional meson exchange interaction with the quark exchange interaction

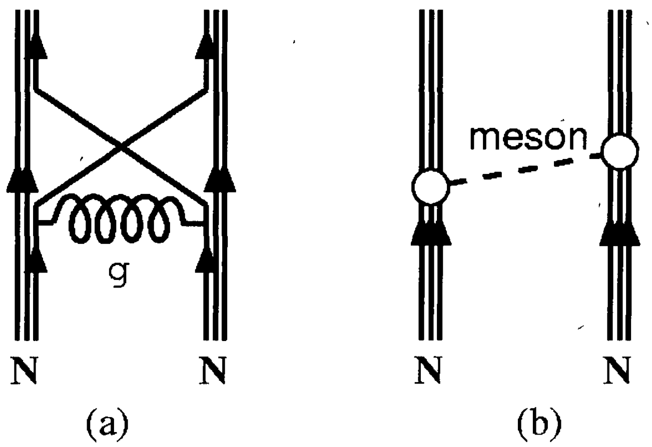

Fig. 3. (a) The quark exchange diagram and (b) the meson exchange diagram for the nucleon-nucleon interaction.

(Fig. 3(b)).[16] The resonating group method is employed in calculating the NN scattering amplitudes employing the quark exchange effect and the model was called "quark cluster model" (QCM). The resulting NN interaction is successful in explaining NN scattering phase shifts and properties of the deuteron. Further elaborations were made by many groups and the NN interaction model based on the short-range repulsion due to the quark exchange mechanism was established.[17]

Because the QCM is based on the quark model, it is straightforward to generalize the model to the strangeness. The first study of the hyperon-nucleon and hyperon-hyperon interactions, including the H dibaryon system, was carried out in 1983.[18] It was found that most of the YN and YY interactions are repulsive at short distances. An exception is the $\Lambda\Lambda - N\Xi - \Sigma\Sigma$ $(I = 0)$ interaction, that has a strong attraction and causes the H dibaryon. The YN and YY interaction model based on the QCM was further developed by the Tokyo group[19] and Kyoto group.[20] Fujiwara gave a summary of recent developments of their model of the NN, YN and YY interactions in detail in the Symposium.

The III was introduced to the QCM first in 1989 in the context of the H dibaryon, which had been predicted to have a too-small mass if the hyperfine interaction is purely the CM interaction. Takeuchi and Oka[8] pointed out that the three-body part of III plays a role to push the H mass higher, so that the H is not a deeply bound state.

It was further pointed out that the III also plays a key role in explaining the spin-orbit interactions among quarks. Both the spectrum of the P wave baryons and the spin-orbit force of the NN scattering are consistently described by the inclusion of the III. The details of the discussion are given in Ref. 21).

§4. Weak decay of strangeness in nuclei

The weak decay of strangeness in nuclear medium provides us with a new opportunity to study hadronic weak interactions. The Λ hypernucleus decays from its ground state by the weak interaction. A free Λ decays into $p\pi^-$ and $n\pi^0$, while new decay modes are allowed in nuclear medium. In the weak decays of heavy hypernuclei, nucleon induced decay modes $\Lambda p \to pn$ and $\Lambda n \to nn$ become dominant. Because these decays do not emit a pion, they are called "nonmesonic weak decay" (NMWD). The reason of the dominance of the NMWD is that the mesonic decay is suppressed by the Pauli exclusion principle, because the momentum of the outgoing nucleon in $\Lambda \to N\pi$ is around 100 MeV/c. On the other hand, in $\Lambda N \to NN$, the outgoing nucleons have larger momenta $\sim$ 400 MeV/c. Recently, these new decay modes are studied in detail in relation to one of the long-standing problem of the hadronic weak interaction, i.e., $\Delta I = 1/2$ rule.

4.1. $\Delta I = 1/2$ *dominance in the nonleptonic strangeness decay*

One of the puzzles in the nonleptonic weak decays of strangeness is the so-called "$\Delta I = 1/2$ rule", i.e., strong enhancement of $\Delta I = 1/2$ amplitudes in strangeness changing weak decays. The experimental ratios of the $\Delta I = 1/2$ and $\Delta I = 3/2$ in the decay amplitudes of kaons and hyperons are about 20, while the standard theory

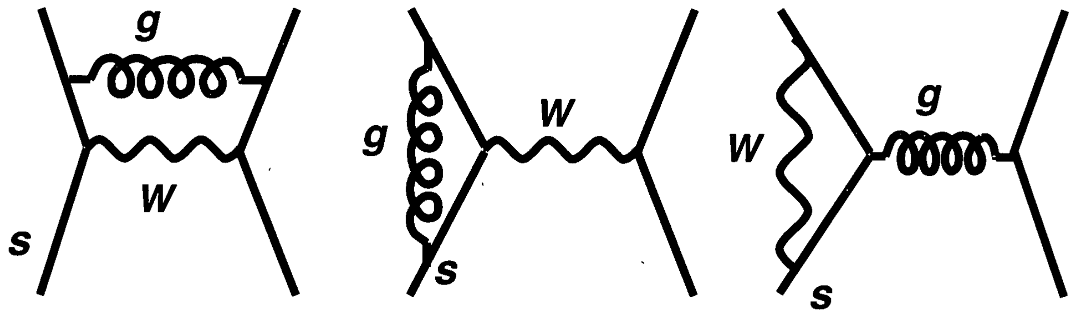

Fig 4 QCD corrections for the $\Delta S = 1$ nonleptonic weak interaction.

of the electroweak interaction predicts that the $\Delta I = 1/2$ and $3/2$ amplitudes are of the same order. Namely, the fundamental vertex given by $s \to u + W^-$, $W^- \to d + \bar{u}$ $(I = 1)$ contains both the $I = 1/2$ and $3/2$ final states.

It is natural to assume that QCD corrections are responsible for the discrepancy. Indeed, the problem is partially solved by considering perturbative QCD corrections on the weak transition process. A standard technique employing a renormalization group improved perturbation theory is shown to enhance (suppress) the $\Delta I = 1/2$ $(\Delta I = 3/2)$ amplitudes.[22] This effect comes from the color-flavor structure of the gluon exchange interactions between quarks (see Fig. 4). The color-magnetic attraction in the scalar diquark, $I = 0$, $S = 0$, $C = \bar{3}$ ud, in the final state enhances the $\Delta I = 1/2$. Also the so-called Penguin diagram (Fig. 4(c)), which is purely $\Delta I = 1/2$ and mixes the right-handed flavor-singlet current, was shown to contribute significantly.

Unfortunately, the perturbative correction is not sufficient to explain the observed $\Delta I = 1/2$ to $\Delta I = 3/2$ ratio. An extra enhancement factor of $3 - 5$ is needed. A possible explanation of the $\Delta I = 1/2$ enhancement in the K decay is $K \to \sigma \to 2\pi$ mechanism.[23] The σ is a scalar $I = 0$ meson of mass ~ 600 MeV. This process enhances the $\Delta I = 1/2$ amplitude because the σ mass is close to the kaon mass.

For hyperon decays, the $\Delta I = 1/2$ enhancement can be explained by a combination of the soft-pion relation and the color symmetry argument. Miura-Minamikawa[24] and Pati-Woo[25] (MMPW) showed that no $\Delta I = 3/2$ amplitude arises from the diagrams in which either the initial two quarks or the final two quarks belong to the same baryon. This theorem is a consequence of the color symmetry of the constituent quarks of the baryon. For pionic decays of hyperons, PV transition matrix elements can be written in terms of baryon-baryon matrix elements by the soft-pion technique, for example,

$$\langle n\pi^0(q)|H^{\rm PV}|\Lambda\rangle \xrightarrow{q\to 0} -\frac{i}{f_\pi}\langle n|[Q_5^0, H^{\rm PV}]|\Lambda\rangle = -\frac{i}{2f_\pi}\langle n|H^{\rm PC}|\Lambda\rangle, \tag{4·1}$$

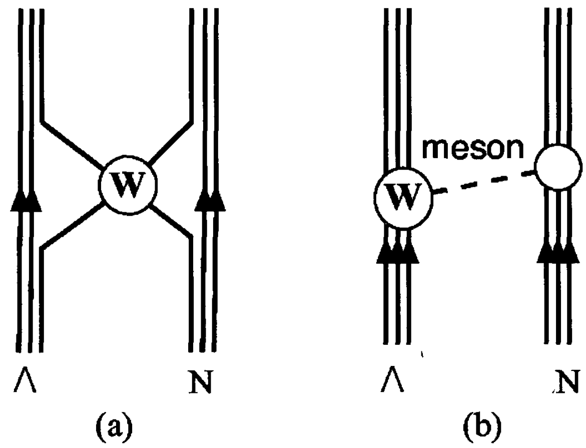

Fig. 5. (a) Direct quark process and (b) meson exchange process for nonmesonic weak decay of Λ hyperon

where we use the following relations satisfied by the weak effective Lagrangian,

$$[Q^a_R, H_W] = 0, \qquad [Q^a_5, H^{\rm PV}] = -[I^a, H^{\rm PC}]. \tag{4·2}$$

Then the baryon-baryon matrix elements, for instance, $\langle n|H^{\rm PC}|\Lambda\rangle$, evaluated in the constituent quark model contain only $\Delta I = 1/2$ transitions due to the MMPW theorem. A similar argument suppresses $\Delta I = 3/2$ amplitudes in the PC transitions, if we assume that pole diagrams are dominant in the PC decays.

Are these mechanisms, which enhance $\Delta I = 1/2$, universal in "all" the nonleptonic weak interactions? In particular, we ask whether the NMWD of hypernuclei satisfies the $\Delta I = 1/2$ rule. In conventional approaches, the nucleon induced weak decays, $\Lambda N \to NN$, are treated as the decay of Λ to $N\pi$ followed by absorption of π with a nucleon in the vicinity (Fig. 5(b)). Because the $\Lambda \to N\pi$ decay is dominated by $\Delta I = 1/2$, the pion exchange process is expected to follow the $\Delta I = 1/2$ rule. However, several recent studies indicated that the $\Delta I = 1/2$ rule may not hold in the NMWD. A possible source of significant $\Delta I = 3/2$ amplitudes is a process called "direct quark" (DQ) transition (Fig. 5(a)).[26),27)] This is a weak interaction process between two baryons, in which a quark in one baryon interacts with a quark in another directly via the effective four-quark weak vertex. Another source pointed out by Shmatikov and Maltman[28)] is a ρ meson exchange diagram, where the weak vertex may contain significant $\Delta I = 3/2$.

Inoue et al.[29)] and Sasaki et al.[30)] derived effective weak transition potentials employing the DQ and π and K exchange processes. They applied the transition potential to hypernuclear decays and calculated the transition amplitudes and decay rates. It was found that the decay rates of $A = 4$ and $A = 5$ hypernuclei are fairly well reproduced. In particular, contributions from the short range weak interactions, DQ and K exchange, are found to solve a long-standing discrepancy between theory and experiment on the ratio of the neutron induced decay Γ_{nn} and the proton induced decay Γ_{pn}, Γ_{nn}/Γ_{pn}.

In order to see the roles of the $\Delta I = 3/2$ transition in the hypernuclear decays, decays of light hypernuclei, $A = 4$ and $A = 5$, are useful.[31)] The ΔI of the weak transition will be clearly seen in the $J=0$ decay amplitudes. We define the ratio, $x=\Gamma^0_{nn}/\Gamma^0_{pn}$, where Γ^0_{pn} (Γ^0_{nn}) denotes the proton induced decay rate in $J=0$. Then x is a clear indicator of ΔI:

$$x=\begin{cases} 2 & \text{for pure } \Delta I=1/2, \\ 1/2 & \text{for pure } \Delta I=3/2. \end{cases} \tag{4·3}$$

The partial decay rates of ${}^4_\Lambda$He, ${}^4_\Lambda$H and ${}^5_\Lambda$He, are the key quantities. We define two ratios of the observables:

$$\alpha \equiv \frac{\Gamma_{NM}({}^4_\Lambda\mathrm{H})}{\Gamma_{NM}({}^4_\Lambda\mathrm{He})}, \qquad \beta \equiv \frac{\Gamma_{nn}({}^5_\Lambda\mathrm{He})}{\Gamma_{pn}({}^5_\Lambda\mathrm{He})}. \tag{4·4}$$

Then the following theorem can be proved if one assumes that the nonmesonic decay is dominated by the two-body processes:

Theorem: If $\alpha > \beta$ then $x < 1/\alpha$.

The current experimental data are not conclusive:

$$\alpha = \frac{0.17\pm 0.11}{0.17\pm 0.05}, \qquad \beta = 0.48 \pm 0.10,$$

but it indicates $x<1$ and therefore possibility of violation of the $\Delta I=1/2$ rule.

A related topic is π^+ decays of hypernuclei. This is another new decay mode that appears only in nuclear medium because $\Lambda \to N\pi$ decay does not emit π^+. A possible decay process is a proton induced transition, $\Lambda p \to nn\pi^+$. We applied the soft pion technique again to the π^+ emission and showed that the decay through $p\Lambda \to n\Sigma^+ \to nn\pi^+$ is strongly hindered in the soft pion (S wave) limit.[32] Instead, the $\Delta I=3/2$ transition gives nonvanishing matrix element in the soft pion limit, and therefore the π^+ emission rate is directly connected to the $\Delta I=3/2$ amplitudes of NMWD. Thus a large π^+ emission rate may indicate the violation of the $\Delta I=1/2$ rule in NMWD.

§5. Conclusion

Studies of hypernuclear physics for these 50 years were very fruitful. Our understanding of the interactions of hyperons and the structures of hypernuclei has reached the level that quantitative predictions are available for comparisons with high quality experimental data.

Yet a new aspect of hypernuclear physics has just begun. We realize that the strangeness plays a key role is studies of low energy phenomena in QCD. The YN and YY interactions are good test ground for the quark model descriptions of the baryon-baryon interactions. In particular, the strange sector of nuclear physics helps to determine the origin of the hyperfine interaction. We stress the importance of the instanton induced interaction (III). The III is indispensable for the pseudoscalar meson spectrum, where the axial $U(1)$ symmetry breaking is observed. In the baryon spectrum, the III gives significant contributions to the mass of the H dibaryon and the spin-orbit interactions in the P-wave baryon excitations.

The nonmesonic weak decays (NMWD) of hypernuclei are fairy well understood in terms of the model with meson exchanges and direct quark (DQ) processes. The long-standing problem of the Γ_{nn}/Γ_{pn} discrepancy between theory and experiment has been solved in cooperation of theoretical and experimental efforts.

The DQ mechanism predicts a large violation of the $\Delta I = 1/2$ rule. The predicted $\Delta I = 1/2$ violation should be tested by experiment. It is crucial to measure $\Gamma_{\rm NM}({}^4_\Lambda{\rm H})$ precisely. The π^+ decay rate may also be useful in determining the role of $\Delta I = 3/2$ component in NMWD.

Strangeness has added a new dimension to the study of low-energy QCD, and has given new approaches to fundamental problems of nonperturbative behavior of low energy QCD. After 50 years, the strangeness is still a key word of new physics. The newly found pentaquark baryon, Θ^+, will lead us to a new era of strangeness hadron physics. A new high-intensity proton accelerator, J-PARC, is under construction and is expected to start new series of hypernuclear experiments in 2007. It will keep experimentalists as well as theorists busy in exploring the world of strangeness further in the 21st century.

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References

1) For recent activities, see also *the Proceedings of the VIII International Conference on Hypernuclear & Strange Particle Physics* (Jefferson Lab, USA, October 14-18, 2003), to be published in Nucl. Phys. A.
2) Y. Akaishi and T. Yamazaki, Phys. Rev. C **65** (2002), 044005.
Y. Akaishi, A. Doté and T. Yamazaki, Prog. Theor. Phys. Suppl. No. 149 (2003), 221.
3) M. Iwasaki et al., nucl-ex/0310018.
4) M. Oka, Prog Theor. Phys. **112** (2004), 1.
5) A DeRujula, H. Georgi and S. L. Glashow, Phys. Rev. D **12** (1975), 147
6) R. L. Jaffe, Phys. Rev. Lett. **38** (1977), 195.
7) For the current status, see T. Sakai, K. Shimizu and K. Yazaki, Prog. Theor. Phys. Suppl. No 137 (2000), 121.
See also M. Oka, Hyperfine Interactions **103** (1996), 275.
8) S. Takeuchi and M Oka, Phys. Rev. Lett. **66** (1991), 1271.
9) G. 't Hooft, Phys. Rev. D **14** (1976), 3432.
10) T. Schäfer and E. Shuryak, Rev. Mod. Phys. **70** (1998), 323.
11) M. Oka and S. Takeuchi, Phys. Rev. Lett. **63** (1989), 1780; Nucl. Phys. A **524** (1991), 649; Phys. Rev. Lett. **66** (1991), 1271.
12) E. Shuryak and J. Rosner, Phys. Lett. B **218** (1989), 72.
13) V. G. Neudatchin, Yu. F. Smirnov and R. Tamagaki, Prog. Theor. Phys. **58** (1977), 1072.
14) M. Oka and K. Yazaki, Phys. Lett. B **90** (1980), 41; Prog Theor. Phys. **66** (1981), 556, ibid **66** (1981), 572.
For a recent review, see Prog. Theor. Phys. Suppl. No. 137 (2000), 1.
15) K. Shimizu, Rep. Prog Phys. **52** (1989), 1.

16) M. Oka and K. Yazaki, Nucl. Phys. A **402** (1983), 477.
17) S. Takeuchi, K Shimizu and K. Yazaki, Nucl. Phys A **504** (1989), 777.
18) M. Oka, K. Shimizu and K. Yazaki, Phys. Lett. B **130** (1983), 365; Nucl. Phys. A **464** (1987), 700.
19) K. Ogawa, S. Takeuchi and M. Oka, in *Properties and Interactions of Hyperons*, ed. B. F. Gibson, P. D. Barnes and K Nakai (World Scientific, 1994), p. 169.
K. Shimizu, S. Takeuchi and A. J. Buchmann, Prog. Theor. Phys. Suppl. No. 137 (2000), 43.
20) Y. Fujiwara, C. Nakamoto and Y. Suzuki, Phys. Rev. Lett. **76** (1996), 2242.
For recent developments, see Y. Fujiwara, Prog. Theor. Phys. Suppl. No. 156 (2004), 17.
21) S. Takeuchi, O. Morimatsu, Y. Tani and M. Oka, Prog. Theor. Phys. Suppl. No. 137 (2000), 83
S. Takeuchi, Phys. Rev. Lett. **73** (1994), 2173, Phys. Rev. D **53** (1996), 6619.
S. Takeuchi, Y. Tani and M. Oka, Nucl. Phys. A **684** (2001), 403.
22) M. K. Gaillard and B. W. Lee, Phys. Rev. Lett. **33** (1974), 108.
G. Altarelli and L. Maiani, Phys. Lett. B **52** (1974), 351.
A. I. Vainshtein, V. I. Zakharov and M. A. Shifman, Sov. Phys. -JETP **45** (1977), 670.
F. J. Gilman and M. B Wise, Phys. Rev. D **20** (1979), 2392.
23) T. Morozumi, C S. Lim and A. I. Sanda, Phys. Rev. Lett. **65** (1990), 404.
T. Inoue, M. Takizawa and M. Oka, Prog. Theor. Phys. Suppl. No 120 (1995), 335.
24) K. Miura and T. Minamikawa, Prog Theor. Phys. **38** (1967), 954.
25) J. C. Pati and C. H. Woo, Phys. Rev. D **3** (1971), 2920.
26) K. Maltman and M. Shmatikov, Phys. Lett B **331** (1994), 1.
27) T. Inoue, S. Takeuchi and M Oka, Nucl. Phys. A **577** (1994), 281.
28) K. Maltman and M. Shmatikov, Phys. Rev. C **51** (1995), 1576.
A. Parreno, A. Ramos, C. Bennhold and K. Maltman, Phys. Lett. B **435** (1998), 1.
29) T Inoue, S. Takeuchi and M. Oka, Nucl. Phys. A **597** (1996), 563.
30) K. Sasaki, T Inoue and M. Oka, Nucl. Phys. A **669** (2000), 331 [Errata; **678** (2000), 455]; ibid. **707** (2002), 477.
31) C Dover, *XIth European Conference on Few-Body Phys.* (1987).
J. Cohen, Phys. Rev. C **42** (1990), 2724.
R A. Schumacher, Nucl Phys A **547** (1992), 143.
32) M Oka, Nucl. Phys. A **647** (1999), 97.