《原子核物理评论》
www.npr.ac.cn Nuclear Physics Review
Started in 1984

基于夸克线图的重子八重态宇称双重态模型

高毕凯　古城徹　原田正康

Parity Doublet Model for Baryon Octets Based on the Quark-line Diagram

GAO Bikai, KOJO Toru, HARADA Masayasu

在线阅读 View online: https://doi.org/10.11804/NuclPhysRev.41.QCS2023.03

引用格式:
高毕凯, 古城徹, 原田正康. 基于夸克线图的重子八重态宇称双重态模型[J]. 原子核物理评论, 2024, 41(3):794–800. doi: 10.11804/NuclPhysRev.41.QCS2023.03
GAO Bikai, KOJO Toru, HARADA Masayasu. Parity Doublet Model for Baryon Octets Based on the Quark–line Diagram[J]. Nuclear Physics Review, 2024, 41(3):794–800. doi: 10.11804/NuclPhysRev.41.QCS2023.03

第 41 卷 第 3 期
2024 年 9 月
原 子 核 物 理 评 论
Nuclear Physics Review
Vol. 41, No. 3
Sep., 2024

Article ID: 1007-4627(2024)03–0794–07

Parity Doublet Model for Baryon Octets Based on the Quark-line Diagram

GAO Bikai[1], KOJO Toru[2], HARADA Masayasu[1,3,4]

(*1. Nagoya University, Nagoya 466-8602, Japan*;
2. Department of Physics, Tohoku University, Sendai 980-8578, Japan;
3. Advanced Science Research Center, Japan Atomic Energy Agency, Tokai 319-1195, Japan;
4. Kobayashi-Maskawa Institute for the Origin of Particles and the Universe, Nagoya 464-8602, Japan)

Abstract: Parity doublet model is an effective chiral model that includes the chiral variant and invariant masses of baryons. The chiral invariant mass has large impacts on the density dependence of models which can be constrained by neutron star observations. In the previous work, models of two-flavors to few times nuclear saturation density have been considered, but in such dense region it is also necessary to consider hyperons. In this work, we construct $\mathrm{SU}(3)_L \otimes \mathrm{SU}(3)_R$ invariant parity doublet models within the linear realization of the chiral symmetry. The major problem in constructing such models has been too many candidates for the chiral representations of baryons. Motivated by the concepts of diquarks and the mended symmetry, we choose the $(3_L, \bar{3}_R) + (\bar{3}_L, 3_R)$, $(3_L, 6_R) + (6_L, 3_R)$ and $(1_L, 8_R) + (8_L, 1_R)$ representations and use quark diagrams to constrain the possible types of Yukawa interactions. The masses of the baryon octets for positive and negative baryons up to the first excitations are successfully reproduced. As expected from the diquark considerations, the ground state baryons are well dominated by $(3_L, \bar{3}_R) + (\bar{3}_L, 3_R)$ and $(1_L, 8_R) + (8_L, 1_R)$ representations, while the excited states require $(3_L, 6_R) + (6_L, 3_R)$ representations. The model should be useful to consider the chiral restoration for strange quarks at large density and the continuity of diquarks from hadronic to quark matter.

Key words: parity doublet; chiral symmetry; quark diagram; neutron star; hyperon

CLC number: O571.6 **Document code:** A

DOI: 10.11804/NuclPhysRev.41.QCS2023.03 **CSTR:** 32260.14.NuclPhysRev.41.QCS2023.03

0 Introduction

Understanding the origin of the nucleon mass is one of the most important subjects in hadron physics. Traditionally, the chiral symmetry in quantum chromodynamics (QCD) and its spontaneous symmetry breaking (SSB)[1−4] are often used to study the nucleon mass. The order parameter of the chiral SSB is the chiral condensate $\langle \bar{q}q \rangle$, which is made of quark-antiquark pairs with different chirality[5−7]. The finite chiral condensate in the vacuum gradually vanishes with increasing density or/and temperature.

The linear sigma model (LSM) is one of effective models broadly used to study the SSB. In the LSM, the order parameter of SSB is the expectation values of the scalar field $\langle \sigma \rangle \propto \langle \bar{q}q \rangle$. In the traditional hadronic model with LSM, the nucleon mass m_N is considered to be generated mainly by chiral condensates. In this case, the nucleon mass vanishes in the high density or temperature region where the chiral symmetry should be restored.

However, there may exist not only the conventional chiral variant mass, but also chiral invariant mass (m_0) whose existence is supported by the previous lattice QCD simulation[8−12]. If so, the nucleon mass would keep a finite value even if the chiral symmetry is restored. The concept of the chiral invariant mass is naturally incorporated in a parity doublet model (PDM) for nucleons in which the ordinary nucleon and its parity partner form a doublet structure[13−20]. Accordingly, nucleons in the PDM are less sensitive to the chiral condensate than in the conventional LSM. This feature affects the construction of nuclear and neutron star (NS) equation of state (EOS), especially the stiffness of the EOS.

In the context of applications to the NS phenomenology, the nuclear EOS within the PDM is often extended to densities beyond the nuclear saturation density $n_0 \simeq 0.16\ \mathrm{fm}^{-3}$. This extrapolation has been achieved in

Received date: 28 Nov. 2023; **Revised date:** 13 Mar. 2024

Biography: GAO Bikai(1997−), male, Jingzhou, Hubei Province, Doctor student, working on nuclear astrophysics; E-mail: gaobikai@hken.phys.nagoya-u.ac.jp

various ways. One approach involves a straightforward extrapolation of the PDM EOS beyond n_0, as studied in Ref. [21]. Another method combines the PDM EOS with a quark model, assuming a quark-hadron crossover, which allows for a smooth transition from hadronic matter to quark matter[22–27]. This hybrid approach, where the PDM EOS is employed up to densities around 2-3 n_0 and interpolate with the quark EOS at $\geqslant 5n_0$ via polynomial interpolants to obtain the unified EOS. This unified EOS is valuable for modeling the behavior of dense matter within NSs and can provide us important insights into the intermediate density region. However, the validity of a purely hadronic picture at densities above $2\,n_0$ is questionable because hyperons, strange baryons containing strange quarks, may enter into the matter and affect its properties. Taking this possibility into consideration, it becomes necessary to explore the effect of strange quark at finite density. In the previous study for the nuclear matter domain, we included the strange quark effects through the Kobayashi-Maskawa-'t Hooft (KMT) interactions and constructed a three-flavor mesonic Lagrangian made of scalar and vector mesons[25]. The corresponding unified EOS was confronted with NS constraints from the NS merger GW170817[28–30], the millisecond pulsar PSR J0030+0451[31], and the maximum mass constraint from the millisecond pulsar PSR J0740+6620[32]. Based on the above NS observation data, we constrained m_0 to the range $400\ \text{MeV} \lesssim m_0 \lesssim 700\ \text{MeV}$ which is more relaxed than in Ref. [23].

In this research, we investigate the influence of strange quarks in the baryonic sector and developed a model based on $\mathrm{SU}(3)_L\times\mathrm{SU}(3)_R$ chiral symmetry. We construct the model in the framework of the parity doublet structure which allow us to introduce the concept of chiral invariant mass. For the study of three flavors case, we begin with chiral representations of quarks with left handed quark $q_L\sim(3_L,1_R)$ and right handed quark $q_R\sim(1_L,3_R)$. With these quark representations, we calculate the possible representations for baryons. Since octet-baryons appear from the representation of $(3_L,\bar{3}_R),(3_L,6_R),(8_L,1_R)$ with $L\leftrightarrow R$, we introduce three types of baryon fields ψ,η and χ respectively. We then include mesonic fields and the parity doublet structure in a chiral invariant way. Both the spectra of baryons with positive and negative parity as well as the first radial excitations are studied. Our previous construction[33] including $(3_L,\bar{3}_R),(8_L,1_R)$ with $L\leftrightarrow R$ representations was motivated by considerations on good diquarks, but did not allow us successful fit to baryon spectra, even after including higher orders of the Yukawa interactions. In this work we manifestly include $(3_L,6_R)+(6_L,3_R)$ representation which is supposed to include bad diquarks. It turns out that manifest inclusion of $(3_L,6_R)+(6_L,3_R)$ allows us to saturate ground state baryons with $(3_L,\bar{3}_R)$ and $(8_L,1_R)$ (with $L\leftrightarrow R$) representations, and now positive and negative parity spectra up to the first excitations can be fitted within physically reasonable model parameters. The description is in line with considerations based on good and bad diquarks.

1 Model construction

The chiral representations of quarks under $\mathrm{SU}(3)_L\times\mathrm{SU}(3)_R$ are written as

$$(q_L)^l\sim(3_L,1_R)\sim(u_L,d_L,s_L)^l, \tag{1}$$

$$(q_R)^l\sim(1_L,3_R)\sim(u_R,d_R,s_R)^r, \tag{2}$$

where the indices $l,\ r=1,2,3$. Since a baryon can be expressed as a direct product of three quarks, we have the following possibilities for the chiral representations of baryons:

$$\begin{aligned}&q_L\otimes(q_L\otimes q_L+q_R\otimes q_R)\sim\\&(1,1)+(8,1)+(8,1)+(10,1)+(3,\bar{3})+(3,6).\end{aligned} \tag{3}$$

The octet baryons are included in $(3,\bar{3}),(3,6)$, and $(8,1)$, which are illustrated in Fig. 1. We introduce the corresponding baryon fields ψ,η,χ and its parity doubling partners $\psi^{\text{mir}},\eta^{\text{mir}},\chi^{\text{mir}}$ as

$$\psi_L\sim(3,\bar{3}),\quad \psi_L^{\text{mir}}\sim(\bar{3},3), \tag{4}$$

$$\eta_L\sim(3,6),\quad \eta_L^{\text{mir}}\sim(6,3), \tag{5}$$

$$\chi_L\sim(8,1),\quad \chi_L^{\text{mir}}\sim(1,8). \tag{6}$$

The right-handed fields are also defined in the similar way. Here, the fields ψ and χ include the $\mathrm{SU}(3)$-flavor octet baryon field $B_{L,R}$ as

$$\psi_{L,R}^{(\text{mrr})}=\frac{1}{\sqrt{3}}\Lambda_{0\,L,R}^{(\text{mrr})}+B_{L,R}^{(\text{mrr})} \tag{7}$$

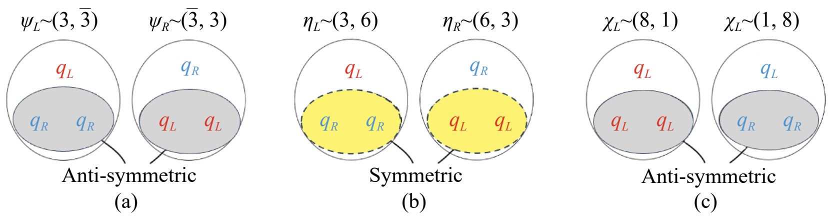

Fig. 1　Quark contents for each baryon representations: (a) $(3,\bar{3})+(\bar{3},3)$, (b) $(3,6)+(6,3)$, and (c) $(8,1)+(1,8)$. The gray shaded diquark indicates flavor antisymmetric representation, while the yellow indicates symmetric one.

$$\chi_{L,R}^{(\mathrm{mrr})} = B_{L,R}^{(\mathrm{mrr})} \tag{8}$$

where Λ_0 is a singlet under the SU(3) flavor symmetry. While the η field as

$$\begin{aligned}(\eta_L\, \eta_R^{\mathrm{mir}})^{(a,\alpha\beta)} &= \frac{1}{\sqrt{6}}\left(\epsilon^{\alpha ac}\delta_k^{\beta} + \epsilon^{\beta ac}\delta_k^{\alpha}\right)(B_L\, B_R^{\mathrm{mir}})_c^k \\ (\eta_R\, \eta_L^{\mathrm{mir}})^{(ab,\alpha)} &= \frac{1}{\sqrt{6}}\left(\epsilon^{\alpha ac}\delta_k^{b} + \epsilon^{bac}\delta_k^{a}\right)(B_R\, B_L^{\mathrm{mir}})_c^k\end{aligned} \tag{9}$$

where $a, b = 1, 2, 3$ are for $\mathrm{SU}(3)_L$ and $\alpha, \beta = 1, 2, 3$ for $\mathrm{SU}(3)_R$. Here the index of ab and $\alpha\beta$ are symmetrized to express "6" representation. Each component of the octet baryon field B is assigned as

$$B = \begin{pmatrix} \frac{1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda & \Sigma^+ & p \\ \Sigma^- & \frac{-1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda & n \\ \Xi^- & \Xi^0 & \frac{-2}{\sqrt{6}}\Lambda \end{pmatrix}. \tag{10}$$

The introduction of the mirror fields enables us to write down the mixing terms as

$$\begin{aligned}\mathcal{L}^{\mathrm{CIM}} =& -m_0^{\psi}\mathrm{tr}\left(\overline{\psi}_{\mathrm{r}}\psi_1^{\mathrm{mir}} - \overline{\psi}_1\psi_{\mathrm{r}}^{\mathrm{mir}}\right) + \mathrm{h.c}- \\ & m_0^{\eta}\left[(\overline{\eta}_r)_{(ab,\alpha)}(\eta_l^{\mathrm{mir}})^{(ab,\alpha)} - (\overline{\eta}_l)_{(a,\alpha\beta)}(\eta_r^{\mathrm{mir}})^{(a,\alpha\beta)}\right] + h.c- \\ & m_0^{\chi}\mathrm{tr}\left(\overline{\chi}_{\mathrm{r}}\chi_1^{\mathrm{mir}} - \overline{\chi}_{\mathrm{l}}\chi_{\mathrm{r}}^{\mathrm{mir}}\right) + \mathrm{h.c}\,,\end{aligned} \tag{11}$$

where the transformation properties, *e.g.*, $\psi_L \to g_L\psi_L g_R^{\dagger}$ and $\psi_R^{\mathrm{mir}} \to g_L\psi_R^{\mathrm{mir}} g_R^{\dagger}$, with $g_{L,R} \in \mathrm{SU}(3)_{L,R}$, make the mass term chiral invariant. The parameter $m_0^{\psi,\eta,\chi}$ corresponds to the chiral invariant mass. Note here that the ψ and χ field contain flavor-antisymmetric diquarks which are called "*good*" diquarks, while (3, 6) contains flavor-symmetric diquark called "*bad*" diquarks. We assume that baryons in representations including "*good*" diquarks are lighter than those including "*bad*" diquarks, so the chiral invariant mass for each baryon field should follow $m_\eta \gtrsim m_\psi \sim m_\chi$. For simplicity, in this research, we set $m_\psi = m_\chi$.

2 Quark diagram for Yukawa interaction

In this section we study the mass spectra of octet-baryons in a model with first-order Yukawa interactions, in which baryons interact with each other with exchange of mesons. For writing the Yukawa interactions, we introduce a 3×3 matrix field M expressing a nonet of scalar and pseudoscalar mesons made of a quark and an anti-quark. The representation under $\mathrm{SU}(3)_L \times \mathrm{SU}(3)_R$ is

$$M \sim (3, \bar{3}). \tag{12}$$

We can then construct the chiral invariant Yukawa interaction terms at the first order of M for several combinations of $\psi, \psi^{\mathrm{mir}}, \eta, \eta^{\mathrm{mir}}, \chi, \chi^{\mathrm{mir}}$ fields. We will give some examples in the following.

For the case only with $(3,\bar{3})+(\bar{3},3)$ representation, the Yukawa interaction couples to a quark q_R in the $(3_L, \bar{3}_R)$ representation. In other words, the matrix M couples to one of quarks forming a good diquark. After the chiral flipping, the $(\bar{3}_L, 3_R)$ representation is formed. The same is true after exchanges of L and R. We then construct the chiral invariant term at the first order in M as

$$\begin{aligned}\mathcal{L}^{\psi} =& g_1\left(\epsilon_{abc}\epsilon^{\alpha\beta\sigma}(\overline{\psi}_R)_\alpha^a(M)_\beta^b(\psi_L)_\sigma^c + h.c\right) + \\ & g_2\left(\epsilon^{abc}\epsilon_{\alpha\beta\sigma}(\overline{\psi}_R^{\mathrm{mir}})_a^\alpha(M)_b^\beta(\psi_L^{\mathrm{mir}})_c^\sigma + h.c\right),\end{aligned} \tag{13}$$

where the ϵ_{ijk} is the totally asymmetric tensor. The corresponding quark diagram is represented in Fig. 2.

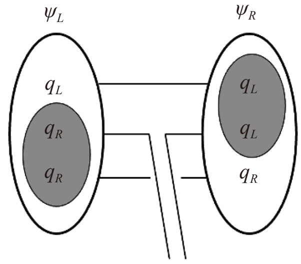

Fig. 2　Yukawa couplings between $(3,\bar{3})$ and $(\bar{3},3)$ baryon fields.

For the Yukawa interactions between ψ and χ, the matrix M couples to a quark q_L in $(\bar{3}_L, 3_R)$ representation forming a "good" diquark. After the chiral flipping the $(3, 6)$ representation is formed. We construct the chiral invariant Lagrangian at the leading order in M as

$$\begin{aligned}\mathcal{L}^{\psi\eta} =& y_1\left[\epsilon_{abc}(\overline{\psi}_R)_\alpha^a(M)_\beta^b(\eta_L)^{(c,\alpha\beta)} + h.c\right] + \\ & y_3\left[\epsilon_{\alpha\beta\sigma}(\overline{\psi}_R^{\mathrm{mir}})_\alpha^a(M)_\beta^b(\eta_L)^{(c,\alpha\beta)} + h.c\right].\end{aligned} \tag{14}$$

The corresponding quark diagram is represented in Fig. 3.

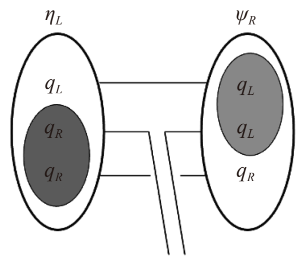

Fig. 3　Yukawa couplings between (3,6) and $(\bar{3},3)$ baryon fields.

For the Yukawa interactions between ψ and χ, a spectator quark q_L in $(3_L, \bar{3}_R)$ representation couples to M and flips the chirality. Then the representation $(1_L, 8_R)$ is formed. We can then construct the following Lagrangian

$$\mathcal{L}^{\psi\chi} = y_2\mathrm{tr}\left(\bar{\psi}_R M^{\dagger}\chi_L + h.c\right) + y_4\mathrm{tr}\left(\bar{\psi}_R^{\mathrm{mir}} M \chi_L^{\mathrm{mir}} + h.c\right) \tag{15}$$

and the corresponding quark diagram as shown in Fig. 4. Similar methods are applied to other types of interactions. Here we have to emphasize that, at the first orders in M, there are no Yukawa interactions that couple χ_L and χ_R fields. This is because the χ field contains three valence quarks with all left-handed or right-handed so that Yukawa interactions with χ should include three quark exchanges that flip the chirality of three quarks.

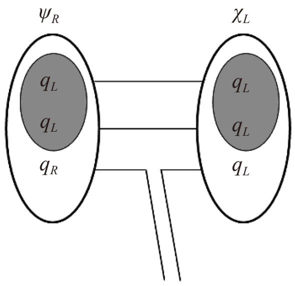

Fig. 4　Yukawa couplings between $(\bar{3},3)$ and $(8,1)$ baryon fields.

3　Numerical fitting

Following Ref. [34], we construct the $\mathrm{SU}(3)_L\times\mathrm{SU}(3)_R$ Lagrangian. We take the mean field approximation in which the meson field M is replaced as its mean field as $\langle M\rangle=\mathrm{diag}(\alpha,\beta,\gamma)$. In the following anaysis we assume the isospin symmetry by taking $\alpha=\beta$. It is convenient to introduce a unified notation for the chiral representations of baryons as $\varPsi_i=(\psi_i,\eta_i,\chi_i,\gamma_5\psi_i^{\mathrm{mir}},\gamma_5\eta_i^{\mathrm{mir}},\gamma_5\chi_i^{\mathrm{mir}})^T$ with $i=N,\varLambda,\varSigma,\varXi$. We then calculate the mass terms of baryons in the form of

$$\mathcal{L}=\sum_{i=N,\varLambda,\varSigma,\varXi}\bar{\varPsi}_i\hat{M}_i\varPsi_i\,, \tag{16}$$

where the $\hat{M}_i(i=N,\varLambda,\varSigma,\varXi)$ is the mass matrices for each baryons. As an example, the mass matrix for nucleon is

$$M_N=\begin{pmatrix} g_1\alpha & -\frac{3y_1}{\sqrt{6}}\alpha & -y_2\alpha & m_0^{\psi} & 0 & 0\\ & \frac{g_3}{2}\alpha & -\frac{3y_5}{\sqrt{6}}\alpha & 0 & m_0^{\eta} & 0\\ & & 0 & 0 & 0 & m_0^{\chi}\\ & & & -g_2\alpha & \frac{3y_3}{\sqrt{6}}\alpha & y_4\alpha\\ & & & & -\frac{g_4}{2}\alpha & \frac{3y_6}{\sqrt{6}}\alpha\\ & & & & & 0 \end{pmatrix}, \tag{17}$$

where we omit the lower triangular part of the matrix M_N, which is understood from the fact that M_N is the symmtric matrix $(M_N)^T=M_N$.

Diagonalizing this 6×6 matrix $\hat{M}_i$, we obtain six mass eigenvalues. We focus on the first four state of baryons in this research as the rest two state are completely predictions with large ambiguity. We determine the vacuum expectation values (VEVs) of the meson field M from the decay constants of pion and kaon as

$$2\alpha=f_\pi,\quad 2\gamma=2f_K-f_\pi\,. \tag{18}$$

In Table 1, the input values of f_π and f_K are shown together with the determined values of α and γ.

Table 1　Physical inputs of the decay constants for pion and kaon[35], and the VEV of the meson field $\langle M\rangle=\mathrm{diag}\{\alpha,\alpha,\gamma\}$.

f_π	93 MeV
f_K	110 MeV
2α	$f_\pi(=93\ \mathrm{MeV})$
2γ	$2f_K-f_\pi(=127\ \mathrm{MeV})$

Using twelve mass values in Table 2, we fit the ten Yukawa coupling constants $g_1,g_2,g_3,g_4,y_1,y_2,y_3,y_4,y_5,y_6$ by minimizing the following function:

$$f_{\min}=\sum_{i=1}^{12}\left(\frac{m_i^{\mathrm{theory}}-m_i^{\mathrm{input}}}{\delta m_i}\right)^2, \tag{19}$$

where errors δm_i are taken as $\delta m_i=10$ MeV for the ground-state baryons and $\delta m_i=100$ MeV for the excited baryons. We take solutions with $f_{\min}-f_{\min}^{\mathrm{best}}<1$ and also put the constrain from the Gell-Mann-Okubo relation by requiring

$$\sum_i|\Delta_{\mathrm{GO,i}}|<100\ \mathrm{MeV}\,, \tag{20}$$

where

$$\Delta_{\mathrm{GO},i}=\frac{m_i[N]+m_i[\varXi]}{2}-\frac{3m_i[\varLambda]+m_i[\varSigma]}{4} \tag{21}$$

with i indicating the octet generations as in Table 2. We show the fitted mass spectrum together with the Gell-Mann-Okubo relation in Fig. 5. In this figure, the blue box shows the experimental values as inputs and the red line shows best fitted values. We find that all the baryon masses below 2 GeV are reproduced well. In previous research in Ref. [33], we use $(3,\bar{3})+(\bar{3},3)$ and (8, 1) + (1, 8) representations to construct a chiral invariant Lagrangian up to second order Yukawa interactions, while the (3, 6) + (6, 3) representations are assumed to be integrated out, leaving some effective interactions. But the mass hierarchy between $\varSigma$ and $\varXi$ cannot be correctly reproduced, indicating the baryon dynamics is not well-saturated by just $(3,\bar{3})+(\bar{3},3)$ and (8, 1) + (1, 8) representations. In this study, on the other hand, we included the (6, 3) + (3, 6) representation and the mass hierarchy between $\varSigma$ and $\varXi$ state is fixed. We also make some prediction about the mass of second excited state of $\varSigma$ and excited states of $\varXi$ below 2 GeV.

We also show the ratio for different representations in the baryon states in the first generation in Fig. 6 where for each combination of (m^{ψ}, m^{η}), the color for each point represent the ratio for $\psi+\chi$ field. The cross marks indicate

Table 2 Physical inputs for the baryon masses belonging to four SU(3)-flavor octets.

J^P	Mass inputs for octet members/MeV			
	N	Λ	Σ	Ξ
$m_1: 1/2^+$(G.S.)	$N(939)$: 939	$\Lambda(1116)$: 1116	$\Sigma(1193)$: 1193	$\Xi(1318)$: 1318
$m_2: 1/2^+$	$N(1440)$: 1440	$\Lambda(1600)$: 1600	$\Sigma(1660)$: 1660	$\Xi(?)$:
$m_3: 1/2^-$	$N(1535)$: 1530	$\Lambda(1670)$: 1674	$\Sigma(?)$:	$\Xi(?)$:
$m_4: 1/2^-$	$N(1650)$: 1650	$\Lambda(1800)$: 1800	$\Sigma(1750)$: 1750	$\Xi(?)$:

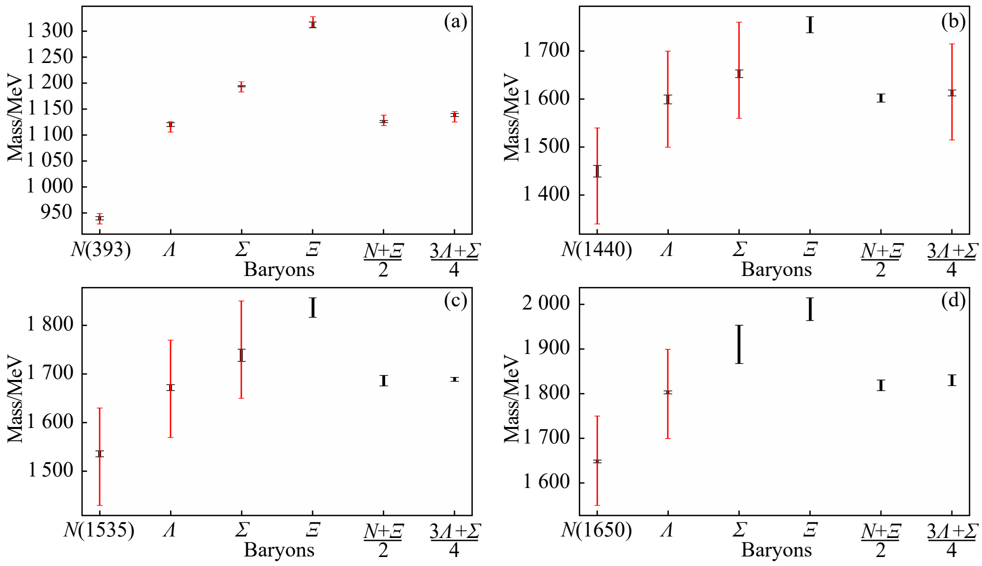

Fig. 5 (color online) Numerical results of the four octet masses for the chiral invariant mass $m^{\psi}=m^{\chi}=700$ MeV and $m^{\eta}=800$ MeV. The blue box shows for the physical inputs with errors and the red line shows for the best fitting value.

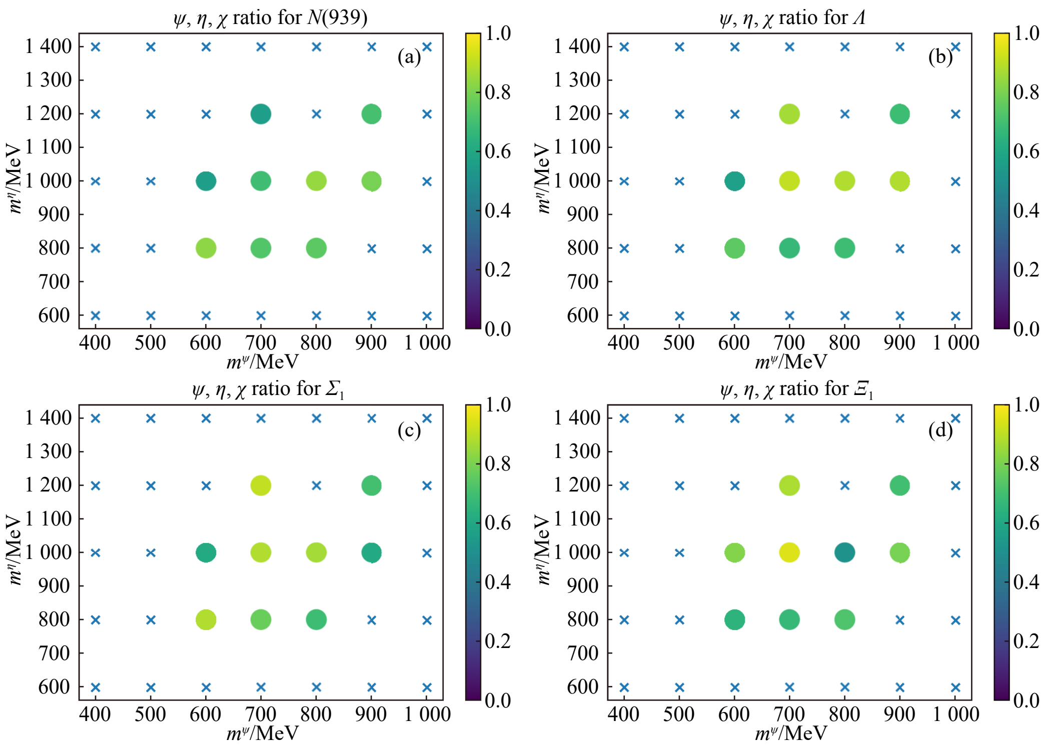

Fig. 6 (color online) Numerical results for the $\psi+\chi$ field ratio for each combination of $(m^{\psi(\chi)}, m^{\eta})$.

that the combinations of (m^{ψ}, m^{η}) are excluded since the combination cannot satisfy $f_{\min} - f_{\min}^{\text{best}} < 1$. From Fig. 6, we find for all the possible combinations of (m^{ψ}, m^{η}), ψ, χ field are dominant in the ground state, which is consistent with the previous research that the inclusion of $(3, \bar{3}) + (3, \bar{3})$ and (8, 1) + (1, 8) is enough to reproduce the ground state baryon masses.

4 Summary and discussion

In this work, we constructed a $\mathrm{SU}(3)_L \times \mathrm{SU}(3)_R$ invariant parity doublet model based on the quark diagram. In the model, mixed state of $(3_L, \bar{3}_R) + (\bar{3}_L, 3_R)$, $(3_L, 6_R) + (6_L, 3_L)$ and $(1_L, 8_R) + (8_L, 1_R)$ representations are included. We use χ^2 fitting of determine 10 parameters in our model to 12 physical inputs and calculate the mass spectrum for the hyperons with mass smaller than 2 GeV. The mass spectrum is well reproduced as shown in Fig. 5. Also we obtained the mixing ratio for different representations. For the ground-state, the results show that, for all the combination of (m^{ψ}, m^{η}), the ψ and χ fields are dominant, indicating the fact that the ground-state is well-saturated by $(3, \bar{3}) + (\bar{3}, 3)$ and (8, 1) + (1, 8) representations. This is consistent with considerations based on diquarks.

References:

[1] WEINBERG S. Phys Rev Lett, 1966, 17: 616.
[2] WEINBERG S. Phys Rev, 1969, 177: 2604.
[3] WEINBERG S. Phys Rev Lett, 1990, 65: 1177.
[4] WEINBERG S. Phys Rev Lett, 2010, 105: 261601.
[5] NAMBU Y, JONA-LASINIO G. Phys Rev, 1961, 122: 345.
[6] NAMBU Y, JONA-LASINIO G. Phys Rev, 1961, 124: 246.
[7] HATSUDA T, KUNIHIRO T. Phys Rept, 1994, 247: 221.
[8] AARTS G, ALLTON C, HANDS S, et al. Phys Rev D, 2015, 92(1): 014503.
[9] AARTS G, ALLTON C, DE BONI D, et al. JHEP, 2017, 06: 034.
[10] AARTS G, ALLTON C, DE BONI D, et al. EPJ Web Conf, 2018, 171: 14005.
[11] AARTS G, ALLTON C, DE BONI D, et al. Phys Rev D, 2019, 99(7): 074503.
[12] AARTS G, ALLTON C, DE BONI D, et al. Springer Proc Phys, 2020, 250: 29.
[13] DETAR C E, KUNIHIRO T. Phys Rev D, 1989, 39: 2805.
[14] JIDO D, NEMOTO Y, OKA M, et al. Nucl Phys A, 2000, 671: 471.
[15] JIDO D, HATSUDA T, KUNIHIRO T. Phys Rev Lett, 2000, 84: 3252.
[16] JIDO D, OKA M, HOSAKA A. Prog Theor Phys, 2001, 106: 873.
[17] NAGATA K, HOSAKA A, DMITRASINOVIC V. Phys Rev Lett, 2008, 101: 092001.
[18] GALLAS S, GIACOSA F, RISCHKE D H. Phys Rev D, 2010, 82: 014004.
[19] GALLAS S, GIACOSA F. Int J Mod Phys A, 2014, 29(17): 1450098.
[20] MOTOHIRO Y, KIM Y, HARADA M. Phys Rev C, 2015, 92(2): 025201.
[21] YAMAZAKI T, HARADA M. Phys Rev C, 2019, 100(2): 025205.
[22] BAYM G, HATSUDA T, KOJO T, et al. Rept Prog Phys, 2018, 81(5): 056902.
[23] MINAMIKAWA T, KOJO T, HARADA M. Phys Rev C, 2021, 103(4): 045205.
[24] MINAMIKAWA T, KOJO T, HARADA M. Phys Rev C, 2021, 104(6): 065201.
[25] GAO B, MINAMIKAWA T, KOJO T, et al. Phys Rev C, 2022, 106(6): 065205.
[26] MINAMIKAWA T, GAO B, KOJO T, et al. Symmetry, 2023, 15(3): 745.
[27] MARCZENKO M, BLASCHKE D, REDLICH K, et al. Universe, 2019, 5(8): 180.
[28] ABBOTT B P, ABBOTT R, ABBOTT T D, et al. Phys Rev Lett, 2017, 119: 161101.
[29] ABBOTT B P, ABBOTT R, ABBOTT T D, et al. Astrophys J Lett, 2017, 848(2): L12.
[30] ABBOTT B P, ABBOTT R, ABBOTT T D, et al. Phys Rev Lett, 2018, 121(16): 161101.
[31] RILEY T E, WATTS A L, RAY P S, et al. Astrophys J Lett, 2021, 918(2): L27.
[32] MILLER M C, LAMB F K, DITTMANN A J, et al. Astrophys J Lett, 2021, 918(2): L28.
[33] MINAMIKAWA T, GAO B, KOJO T, et al. Phys Rev D, 2023, 108(7): 076017.
[34] NISHIHARA H, HARADA M. Phys Rev D, 2015, 92(5): 054022.
[35] WORKMAN R L, BURKERT V D, CREDE V, et al. PTEP, 2022, 2022: 083C.

基于夸克线图的重子八重态宇称双重态模型

高毕凯[1,1)], 古城徹[2], 原田正康[1,3,4]

(1. 名古屋大学, 日本 名古屋 466-8602；
2. 日本东北大学, 日本 仙台 980-8578；
3. 素粒子宇宙起源研究所, 日本 名古屋 319-1195；
4. 日本原子能研究开发机构, 日本 东海 464-8602)

摘要: 宇称双重态模型是一种有效的手征模型，包含了重子的手征变量质量和手征不变质量。手征不变质量对模型的密度依赖性有很大影响，可以通过中子星的观测来约束。在之前的工作中，已经考虑了核饱和密度几倍范围内的两味模型，但在如此高的密度下，也有必要考虑超子。在本工作中，在手征对称性的线性框架内构建了 $SU(3)_L \otimes SU(3)_R$ 不变的宇称双重态模型。构建这类模型的主要问题是重子的手征表示候选太多。受到双夸克和手征对称性概念的启发，选择了 $(3_L, \bar{3}_R) + (\bar{3}_L, 3_R)$、$(3_L, 6_R) + (6_L, 3_R)$ 和 $(1_L, 8_R) + (8_L, 1_R)$ 的重子手征表示，并使用夸克图来约束可能的Yukawa相互作用类型。成功地重现了正宇称和负宇称重子八重态的质量，直至第一激发态。正如双夸克考虑所预期的那样，基态重子主要由 $(3_L, \bar{3}_R) + (\bar{3}_L, 3_R)$ 和 $(1_L, 8_R) + (8_L, 1_R)$ 表示主导，而激发态则需要包含 $(3_L, 6_R) + (6_L, 3_R)$ 来解释。该模型对于考虑高密度下奇异夸克的手征恢复以及从强子物质到夸克物质的双夸克连续性提供了基础框架。

关键词: 宇称双重态；手征对称性；夸克图；中子星；超子

收稿日期: 2023-11-28； **修改日期:** 2024-03-13

1) E-mail: gaobikai@hken.phys.nagoya-u.ac.jp