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Axial charges of N(1535) and N(1650) in two-flavor lattice QCD*

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Abstract We show the lattice QCD results on the axial charge $g_A^{N^*N^*}$ of negative parity nucleon resonances, $N^*(1535)$ and $N^*(1650)$, which are key clues to the chiral structure in baryon sector. The measurements are performed with up and down dynamical quarks employing the renormalization-group improved gauge action at β=1.95 and the $\mathcal{O}(a)$ improved clover quark action with the hopping parameters, κ=0.1375, 0.1390 and 0.1400. In order to properly separate signals of $N^*(1535)$ and $N^*(1650)$, we construct 2×2 correlation matrices and diagonalize them. Wraparound contributions in the correlator, which can be another source of signal contaminations, are eliminated by imposing the Dirichlet boundary condition in the temporal direction. We find that the axial charge of $N^*(1535)$ takes small values as $g_A^{N^*N^*}\sim\mathcal{O}(0.1)$, whereas that of $N^*(1650)$ is about 0.5, which is found independent of quark masses and consistent with the predictions by the naive nonrelativistic quark model.

Key words axial charge, chiral symmetry, negative parity nucleon

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1 Introduction

Quantum ChromoDynamics (QCD) is widely known as the underlying theory of hadron or nuclear physics. Due to its strong coupling nature, it still remains difficult to describe hadron or nuclear systems directly in terms of quarks' and gluons' degrees of freedom. Nevertheless, the properties of QCD are reflected on several fronts in hadron physics, and frequently serve as important clues as to clarification of hadron dynamics. One of the most important and interesting concepts in QCD is chiral symmetry. Chiral symmetry is an approximate global symmetry in QCD, and this symmetry together with its spontaneous breaking has been one of the key ingredients in the low-energy hadron or nuclear physics. While up and down quarks originally have current quark masses of the order of a few MeV, the spontaneous breaking of chiral symmetry gives rise to the large constituent quark masses of a few hundreds MeV. The spontaneous symmetry breaking is in this sense responsible for about 99% of mass of the nucleon and hence that of our world. One may thus consider that chiral condensate $\langle\bar{\psi}\psi\rangle$, the order parameter of the chiral phase transition, plays an essential role in the hadron-mass genesis in the light quark sector. On the other hand, chiral symmetry is restored in systems where hard external energy scales such as high-momentum transfer, temperature (T), baryon density and so on exist, owing to the asymptotic freedom of QCD. Then, several questions arise: Are all hadronic modes massless in such systems? Can hadrons be massive even without non-vanishing chiral condensate?

An interesting possibility was suggested by DeTar and Kunihiro[1]. Motivated by the lattice QCD results[1—3] which reported finite-mass parity-doubled baryons at finite temperature, they constructed a model in which nucleons can be massive even without the help of chiral condensate introducing a possible chirally invariant mass term. This new type of mass

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term gives degenerated nonvanishing masses to the members in a chiral multiplet (a nucleon and its parity partner in their example) even when chiral condensate is set to zero. To demonstrate this possibility for a finite-T case, they introduced a linear sigma model which offers a nontrivial chiral structure in the baryon sector and a mass-generation mechanism essentially different from that by the spontaneous chiral symmetry breaking. Interestingly enough, their chiral doublet model has recently become a source of debate as a possible scenario of observed parity doubling in excited baryons[4—9], although their work[1] was originally applied to finite-T systems.

It is thus an interesting problem to clarify the chiral structure of excited baryons in the light quark sector beyond model considerations. One of the key observables sensitive to the chiral structure of the baryon sector is axial charges[1]. The axial charge of a nucleon N is encoded in the three-point function

$$\langle \mathrm{N}|A_\mu^a|\mathrm{N}\rangle = \bar{u}\frac{\tau^a}{2}[\gamma_\mu\gamma_5 g_A(q^2)+q_\mu\gamma_5 h_A(q^2)]u. \quad (1)$$

Here, $A_\mu^a \equiv \bar{Q}\gamma_\mu\gamma_5\frac{\tau^a}{2}Q$ is the isovector axial current, and q denotes a transferred momentum. The axial charge g_A is defined by $g_A(q^2)$ at $q^2=0$. It is a well-know fact that the axial charge g_A^{NN} of N(940) is 1.26. Though the axial charges in the chiral broken phase can be freely adjusted with higher-dimensional possible terms and cannot be the crucial clues for the chiral structure[5, 6], they would surely reflect the internal structure of baryons and would play an important role in the clarification of the low-energy hadron dynamics.

We show the first unquenched lattice QCD study[10] of the axial charges $g_A^{\mathrm{N^*N^*}}$ of N*(1535) and N*(1650). We adopt $16^3\times 32$ gauge configurations with two flavors of dynamical quarks, generated by CP-PACS collaboration[11] with the renormalization-group improved gauge action and the mean-field improved clover quark action. We choose the gauge configurations at $\beta = 1.95$ with the clover coefficient $c_{\mathrm{SW}} = 1.530$, whose lattice spacing a is determined as 0.1555(17) fm[11] so that the ρ meson mass is reproduced. We perform measurements with 590, 680, and 680 gauge configurations with three different hopping parameters for sea and valence quarks, $\kappa_{\mathrm{sea}},\kappa_{\mathrm{val}} = 0.1375, 0.1390$ and 0.1400, which correspond to quark masses of $\sim$ 150, 100, 65 MeV and the related π-ρ mass ratios are $m_{\mathrm{PS}}/m_{\mathrm{V}} = 0.804(1)$, 0.752(1) and 0.690(1), respectively. Statistical errors are estimated by the jackknife method with the bin size of 10 configurations.

Our main concern is the axial charges of the negative-parity nucleon resonances N*(1535) and N*(1650) in $\frac{1}{2}^-$ channel. We then have to construct an optimal operator which dominantly couples to N*(1535) or N*(1650). We employ the following two independent nucleon fields,

$$N_1(x) \equiv \varepsilon_{\mathrm{abc}} u^a(x)(u^b(x)C\gamma_5 d^c(x)) \quad (2)$$

and

$$N_2(x) \equiv \varepsilon_{\mathrm{abc}} \gamma_5 u^a(x)(u^b(x)Cd^c(x)), \quad (3)$$

in order to separate signals of N*(1535) and N*(1650) with variational analyses. (Here, $u(x)$ and $d(x)$ are Dirac spinor for u- and d- quark, respectively, and a, b, c denote the color indices.) Even after the successful signal separations, there still remain several signal contaminations mainly because lattices employed in actual calculations are finite systems: Signal contaminations by scattering states and those by wraparound effects.

Since our gauge configurations are unquenched ones, the negative parity nucleon states could decay to π and N, and their scattering states could come into the spectrum. However, the sum of the pion mass M_π and the nucleon mass M_{N} is in our setups heavier than the masses of the lowest two states (would-be N*(1535) and N*(1650)) in the negative parity channel. We then do not suffer from any scattering-state signals.

The other possible contamination is wraparound effects[12]. Let us consider a two-point baryonic correlator $\langle \mathrm{N}^*(t_{\mathrm{snk}})\bar{\mathrm{N}}^*(t_{\mathrm{src}})\rangle$ in a Euclidean space-time. Here, the operators $\mathrm{N}^*(t)$ and $\bar{\mathrm{N}}^*(t)$ have nonzero matrix elements, $\langle 0|\mathrm{N}^*(t)|\mathrm{N}^*\rangle$ and $\langle \mathrm{N}^*|\bar{\mathrm{N}}^*(t)|0\rangle$, and couple to the state $|\mathrm{N}^*\rangle$. Since we perform unquenched calculations, the excited nucleon N* can decay into N and π, and even when we have no scattering state $|\mathrm{N}+\pi\rangle$, we could have another "scattering states". The correlator $\langle \mathrm{N}^*(t_{\mathrm{snk}})\bar{\mathrm{N}}^*(t_{\mathrm{src}})\rangle$ can still accommodate, for example, the following term.

$$\langle \pi|\mathrm{N}^*(t_{\mathrm{snk}})|\mathrm{N}\rangle\langle \mathrm{N}|\bar{\mathrm{N}}^*(t_{\mathrm{src}})|\pi\rangle \times \mathrm{e}^{-E_{\mathrm{N}}(t_{\mathrm{snk}}-t_{\mathrm{src}})}\times \mathrm{e}^{-E_\pi(N_t-t_{\mathrm{snk}}+t_{\mathrm{src}})}. \quad (4)$$

Here, N_t denotes the temporal extent of a lattice. Such a term is quite problematic and mimic a fake plateau at $E_{\mathrm{N}}-E_\pi$ in the effective mass plot because it behaves as $\sim \mathrm{e}^{-(E_{\mathrm{N}}-E_\pi)(t_{\mathrm{snk}}-t_{\mathrm{src}})}$. Although these contaminations disappear when one employ enough large-N_t lattice, our lattices do not have so large N_t. In order to eliminate such contributions, we impose the Dirichlet condition on the temporal boundary for valence quarks, which prevents valence quarks from going over the boundary. Though the boundary is still transparent for the states with the same quantum numbers as vacuum, e.g. glueballs, such contributions will be suppressed by the factor of $\mathrm{e}^{-E_{\mathrm{G}}N_t}$ and we neglect them in this paper. We note here that wraparound effects can be found even in quenched calculations[12].

2 Formulation

We give a brief introduction to our formulation[12, 13]. Let us assume that we have a set of N independent operators, O^I_{snk} for sinks and $O^{I\dagger}_{\mathrm{src}}$ for sources. We can then construct an $N\times N$ correlation matrix

$$\mathcal{C}^{IJ}(T)\equiv\langle O^I_{\mathrm{snk}}(T)O^{J\dagger}_{\mathrm{src}}(0)\rangle = C^\dagger_{\mathrm{snk}}\Lambda(T)C_{\mathrm{src}}\,. \tag{5}$$

Here,

$$(C^\dagger_{\mathrm{snk}})_{Ii}\equiv\langle\mathrm{vac}|O^I_{\mathrm{snk}}|i\rangle \tag{6}$$

and

$$(C_{\mathrm{src}})_{jI}\equiv\langle j|O^{J\dagger}_{\mathrm{src}}|\mathrm{vac}\rangle \tag{7}$$

are general matrices, and

$$\Lambda(T)_{ij} \tag{8}$$

is a diagonal matrix given by

$$\Lambda(T)_{ij}\equiv\delta_{ij}\mathrm{e}^{-E_iT}. \tag{9}$$

The optimal source and sink operators, $\mathcal{O}^{i\dagger}_{\mathrm{src}}$ and $\mathcal{O}^i_{\mathrm{snk}}$, which couple dominantly (solely in the ideal case) to i-th lowest state, are obtained as

$$\mathcal{O}^{i\dagger}_{\mathrm{src}}=\sum_J O^{J\dagger}_{\mathrm{src}}(C_{\mathrm{src}})^{-1}_{Ji} \tag{10}$$

and

$$\mathcal{O}^i_{\mathrm{snk}}=\sum_J (C^\dagger_{\mathrm{snk}})^{-1}_{iJ}O^J_{\mathrm{snk}}, \tag{11}$$

since

$$(C^\dagger_{\mathrm{snk}})^{-1}\mathcal{C}(T)(C_{\mathrm{src}})^{-1}=\Lambda(T) \tag{12}$$

is diagonal. Besides overall constants, $(C_{\mathrm{src}})^{-1}$ and $(C^\dagger_{\mathrm{snk}})^{-1}$ are obtained as the right and left eigenvectors of $\mathcal{C}^{-1}(T+1)\mathcal{C}(T)$ and $\mathcal{C}(T)\mathcal{C}(T+1)^{-1}$, respectively.

The zero-momentum-projected point-type operators,

$$N_1(t)\equiv\sum_{\boldsymbol{x}}\varepsilon_{\mathrm{abc}}u^a(\boldsymbol{x},t)(u^b(\boldsymbol{x},t)C\gamma_5d^c(\boldsymbol{x},t)) \tag{13}$$

and

$$N_2(t)\equiv\sum_{\boldsymbol{x}}\varepsilon_{\mathrm{abc}}\gamma_5u^a(\boldsymbol{x},t)(u^b(\boldsymbol{x},t)Cd^c(\boldsymbol{x},t)), \tag{14}$$

are chosen for the sinks. For the sources, we employ the following wall-type operators in the Coulomb gauge,

$$\overline{N_1}(t)\equiv\sum_{\boldsymbol{x}_1,\boldsymbol{x}_2,\boldsymbol{x}_3}\varepsilon_{\mathrm{abc}}\bar{u}^a(\boldsymbol{x}_1,t)(\bar{u}^b(\boldsymbol{x}_2,t)C\gamma_5\bar{d}^c(\boldsymbol{x}_3,t)) \tag{15}$$

and

$$\overline{N_2}(t)\equiv\sum_{\boldsymbol{x}_1,\boldsymbol{x}_2,\boldsymbol{x}_3}\varepsilon_{\mathrm{abc}}\gamma_5\bar{u}^a(\boldsymbol{x}_1,t)(\bar{u}^b(\boldsymbol{x}_2,t)C\bar{d}^c(\boldsymbol{x}_3,t)). \tag{16}$$

The parity is flipped by multiplying the operator by γ_5; $N_i^+(t)\equiv N_i(t)$ and $N_i^-(t)\equiv\gamma_5N_i(t)$. The optimized sink (source) operators $\mathcal{N}_i^\pm$ ($\overline{\mathcal{N}_i^\pm}$), which couple dominantly to the i-th lowest state are constructed as

$$\begin{aligned}\mathcal{N}_i^\pm(t) &= N_1^\pm(t)+\left[(C_{\mathrm{snk}}^{\pm\dagger})^{-1}_{i2}/(C_{\mathrm{snk}}^{\pm\dagger})^{-1}_{i1}\right]N_2^\pm(t)\equiv\\ &\quad N_1^\pm(t)+L_i^\pm N_2^\pm(t),\end{aligned} \tag{17}$$

and

$$\begin{aligned}\overline{\mathcal{N}_i^\pm}(t) &= \overline{N_1^\pm}(t)+[(C_{\mathrm{src}}^{\pm})^{-1}_{2i}/(C_{\mathrm{src}}^{\pm})^{-1}_{1i}]\overline{N_2^\pm}(t)\equiv\\ &\quad \overline{N_1^\pm}(t)+R_i^\pm\overline{N_2^\pm}(t).\end{aligned} \tag{18}$$

With such optimized operators, we can easily compute the (non-renormalized) vector and axial charges $g^{\pm[\mathrm{lat}]}_{V,A}$ for the positive- and negative-parity nucleons via three-point functions with the so-called sequential-source method[14]. In practice, we evaluate $g^{\pm[\mathrm{lat}]}_{V,A}(t)$ defined as

$$g^{\pm[\mathrm{lat}]}_{V,A}(t)=\frac{\mathrm{Tr}\ \Gamma_{V,A}\langle B(t_{\mathrm{snk}})J^{V,A}_\mu(t)\overline{B}(t_{\mathrm{src}})\rangle}{\mathrm{Tr}\ \Gamma_{V,A}\langle B(t_{\mathrm{snk}})\overline{B}(t_{\mathrm{src}})\rangle}\,, \tag{19}$$

and extract $g^{\pm[\mathrm{lat}]}_{V,A}$ by the fit $g^{\pm[\mathrm{lat}]}_{V,A}=g^{\pm[\mathrm{lat}]}_{V,A}(t)$ in the plateau region. Here, $B(t)$ denotes the (optimized) interpolating field for nucleons, and $\Gamma_{V,A}$ are $\gamma_\mu\frac{1+\gamma_4}{2}$ and $\gamma_\mu\gamma_5\frac{1+\gamma_4}{2}$, respectively. $J^{V,A}_\mu(t)$ are the vector and the axial vector currents inserted at t. We show in Fig. 1 $g_A^{-0[\mathrm{lat}]}(t)$ for N*(1535) as a function of the current insertion time t. They are rather stable around $t_{\mathrm{src}}<t<t_{\mathrm{snk}}$.

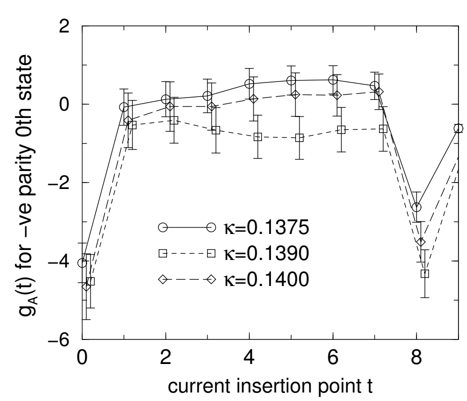

Fig. 1. The non-renormalized axial charge of N*(1535), $g_A^{-0[\mathrm{lat}]}(t)$, as a function of the current insertion time t.

We finally construct the renormalized charges $g^{\pm}_{A,V} = \widetilde{Z}_{A,V} g^{\pm[\text{lat}]}_{A,V}$ with the prefactors $\widetilde{Z}_{A,V} \equiv 2\kappa u_0 Z_{V,A}\left(1+b_{V,A}\dfrac{m}{u_0}\right)$, which can be estimated with the values listed in Ref. [11].

3 Results

We first take a look at the vector charges $g_V^{0\pm}$ of the ground-state positive- and negative-parity nucleons as well as the axial charge g_A^{0+} of the ground-state positive-parity nucleon, which are well known and can be the references. We show $g_V^{0\pm}$, the vector charges of the positive- and the negative-parity nucleons at κ=0.1375, 0.1390 and 0.1400, in the upper panel in Fig. 2, where the vertical axis denotes $g_V^{0\pm}$ and the horizontal one the squared pion masses. The vector charges should be unity if they are conserved, whereas we can actually find about 10% deviations (See the upper panel in Fig. 2). Such deviations are considered to arise due to the discretization errors in perturbative renormalization factors: The present lattice spacing is about 0.15 fm, which is far from the continuum limit. In fact, the decay constants obtained with the same setup as ours deviate from the continuum values by $\mathcal{O}(10)$%. We should then count at least 10% ambiguities in our results. The axial charge g_A^{0+} of the positive parity nucleon is also shown in the lower panel in Fig. 2. As found in the previous lattice studies, the axial charge of the positive parity nucleon shows little quark-mass dependence, and they lie around the experimental value 1.26.

We finally show the axial charges of the negative-parity nucleon resonances in the lower panel in Fig. 2. One finds at a glance that the axial charge g_A^{0-} of N*(1535) takes quite small value as $g_A^{0-} \sim \mathcal{O}(0.1)$, and that even the sign is quark-mass dependent. While the wavy behavior might come from the sensitiveness of g_A^{0-} to quark masses, this behavior may indicate that g_A^{0-} is rather consistent with zero. These small values are not the consequence of the cancellation between u- and d-quark contributions. The u- and d-quark contributions to g_A^{0-} are in fact individually small[10]. We additionally make some trials with lighter u- and d-quark masses at κ=0.1410. Since we have less gauge configurations and the statistical fluctuation is larger at this kappa, we fail to find a clear plateau in the effective mass plots of the two-point correlators and the extracted masses of the negative-parity states cannot be very reliable. Leaving aside these failures, we try to extract g_A^{0-}. The result is added in the lower panel in Fig. 2 as a faint-colored symbol, which is consistent with those obtained at other κ's.

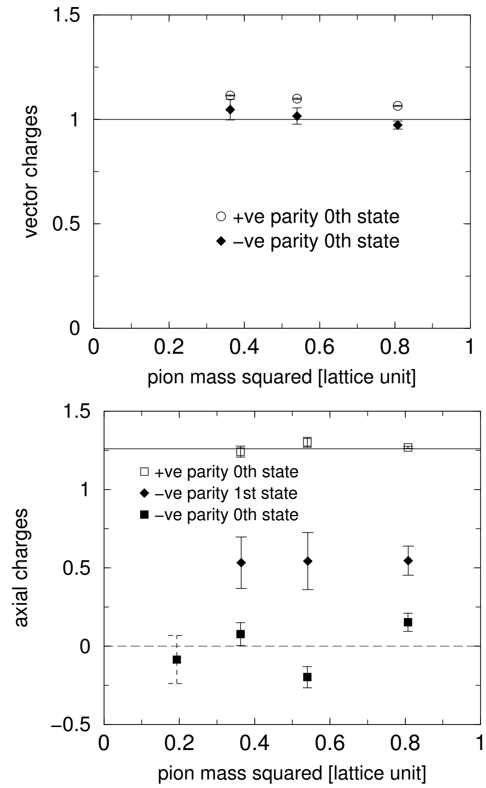

Fig. 2. The renormalized vector and axial charges of the positive- and the negative-parity nucleons are plotted as a function of the squared pion mass m_π^2. upper panel: The results of the vector charges. The solid line is drawn at $g_V = 1$ for reference. lower panel: The results of the axial charges. The solid line is drawn at $g_A = 1.26$ and the dashed line is drawn at $g_A = 0$.

On the other hand, the axial charge g_A^{1-} of N*(1650) is found to be about 0.55, which has almost no quark-mass dependence. The striking feature is that these axial charges, $g_A^{0-} \sim 0$ and $g_A^{1-} \sim 0.55$, are consistent with the naive nonrelativistic quark model calculations[15, 16], $g_A^{0-} = -\dfrac{1}{9}$ and $g_A^{1-} = \dfrac{5}{9}$. These values are obtained assuming that the wave functions of N*(1535) and N*(1650) are $|l=1, S=\dfrac{1}{2}\rangle$ and $|l=1, S=\dfrac{3}{2}\rangle$ neglecting the possible state mixing. (Here, l denotes the orbital angular momentum and S the total spin.)

4 Discussions

In the chiral doublet model[1, 4], the small $g_A^{\mathrm{N^*N^*}}$ is realized when the system is decoupled from the chiral condensate $\langle\bar{\psi}\psi\rangle$. The small g_A^{0-} of N*(1535) then does not contradict with the possible and attempting scenario, chiral restoration scenario in excited hadrons[4]. If this scenario is the case, the origin of mass of N*(1535) (or excited nucleons) is essentially different from that of the positive-parity ground-state nucleon N(940), which mainly arises from the spontaneous chiral symmetry breaking. However, the non-vanishing axial charge of N*(1650) unfortunately gives rise to doubts about the scenario.

In order to reveal the realistic chiral structure, studies with much lighter u,d quarks will be indispensable, since the present quark masses are rather heavy. A study of the axial charge of Roper as well as the inclusion of strange sea quarks could also cast light on the low-energy chiral structure of baryons and the origin of mass. Especially, the large $\bar{s}s$ component in N(1535) was pointed out[17], and hence the inclusion of dynamical strange quark can play a crucial role.

Negative parity baryons other than nucleons are also interesting particles that give us the knowledge about hadron dynamics. For example, Λ(1405) is attracting much interest in terms of its multiquark components, where important dynamics including strange quarks is well encoded[18—21]. Lattice QCD calculations would be helpful to reveal the properties of such particles[22] as well as excited nucleon resonances.

5 Conclusions

We have reported the first lattice QCD study of the axial charge $g_A^{N^*N^*}$ of N*(1535) and N*(1650), with two flavors of dynamical quarks. We employed the renormalization-group improved gauge action at β=1.95 and the mean-field improved clover quark action with the hopping parameters, κ=0.1375, 0.1390 and 0.1400. We have found the small axial charge g_A^{0-} of N*(1535), whose absolute value seems less than 0.2 and which is almost independent of quark mass, whereas the axial charge g_A^{1-} of N*(1650) is found to be about 0.55. These values are consistent with the naive nonrelativistic quark model predictions, and could not be the favorable evidences for the chiral restoration scenario[4] in low-lying excited hadrons. Further investigations on the axial charges of N*(1535) or other excited baryons will cast light on the chiral structure of the low-energy hadron dynamics and on where hadronic masses come from.

All the numerical calculations were performed on NEC SX-8R at RCNP and CMC, Osaka University, on SX-8 at YITP, Kyoto University, and on BlueGene at KEK. The unquenched gauge configurations employed in our analysis were all generated by CP-PACS collaboration[11].

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