Suppression of πNN^* Coupling and Chiral Symmetry

D. Jido* and M. Oka
Department of Physics, Tokyo Institute of Technology, Meguro, Tokyo 152 Japan

A. Hosaka
Numazu College of Technology, 3600 Ooka, Numazu 410 Japan
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Meson-baryon couplings between positive and negative parity baryons are investigated using two-point correlation functions in the soft meson limit. We find that the πNN^* coupling vanishes due to chiral symmetry, while the ηNN^* coupling remains finite. We perform an analysis based on the algebraic method for SU (2) and SU (3) chiral symmetry, and find that baryon axial charges play an essential role for vanishing coupling constants. [S0031-9007(97)05052-7]

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Chiral symmetry, with its spontaneous breakdown in the ground state, is one of the most important symmetries in QCD. Current algebra and low energy theorems have been successfully applied to hadronic phenomena such as meson-meson scattering and photoproduction of mesons [1]. Recent interest in chiral symmetry is largely related to the study of the restoration of chiral symmetry at finite temperature and density [2]. One of the implications there is the appearance of degenerate particles with opposite parties when the chiral symmetry is restored. For the meson sector they would be, for example, $(\sigma, \vec{\pi})$ and $(\vec{\rho}, \vec{a}_1)$. In contrast, the role of chiral symmetry is not fully explored in the baryon sector. Possible candidates for such would-be degenerate particles are $N(939)$ and $N(1535)$. However, the fate of these particles under chiral symmetry restoration is not clear.

Recently, Jido, Kodama, and Oka [3] studied negative parity baryons B^* in the QCD sum rule [4,5]. There a new method of extracting the information of B^* from correlation functions has been formulated and various baryon masses were computed. The resulting masses turned out to be in good agreement with data. It was then pointed out that the mass splittings between positive and negative parity baryons are produced by a nonzero quark condensate, suggesting that chiral symmetry is playing an important role in the negative parity baryons.

We here study meson-baryon couplings as a quantity which reflects chiral symmetry. Experimentally the decay modes of $N(1535)$ are well known, and using the observed decay widths [6], we find two major coupling constants, $g_{\pi NN^*} \sim 0.7$ and $g_{\eta NN^*} \sim 2$. Here we realize that the values of $g_{\pi NN^*} \sim 0.7$ is much smaller than the typical values of pion-baryon coupling constants, e.g., $g_{\pi NN} \sim 13$. The purpose of the present paper is to show that the chiral property of the baryon plays a crucial role in suppressing $g_{\pi NN}$.

Let us start with the interpolating field for the nucleon. Without derivatives, it is given by the superposition of two independent terms [7]:

$$J(x;t) = \epsilon^{abc}\{[u_a(x)Cd_b(x)]\gamma_5 u_c(x) + t[u_a(x)C\gamma_5 d_b(x)u_c(x)\}, \tag{1}$$

where a, b, and c are color indices, and $C = i\gamma_2\gamma_0$ is the charge conjugation matrix. The parameter t controls the coupling strength of J to various nucleons. For instance, it is known that Ioffe's current $J(x; t = -1)$ [8] couples strongly to the positive parity nucleon [7], while it was found in Ref. [3] that $J(x; t = 0.8)$ is optimal for the negative parity nucleon. In this way, the interpolating field (1) has been successfully used in the QCD sum rule to study both positive and negative parity nucleons [3]. Later we will see that the parameter t also specifies the chiral structure of $J(x; t)$, which determines the property of meson-nucleon couplings.

Now let us investigate meson-nucleon coupling constants. We consider the πNN^* coupling, where N^* denotes the negative parity nucleon. Following the method used by Shiomi and Hatsuda [9], we study two-point correlation functions between the vacuum and a one-pion state. In the soft pion limit ($q^\mu \to 0$), the correlation function is given by

$$\Pi^\pi(p) = i\int d^4x\, e^{ipx}\langle 0|TJ(x;s)\bar{J}(0;t)|\pi(q\to 0)\rangle = i\Big[\Pi_0^\pi(p^2)\gamma_5 + \Pi_1^\pi(p^2)\not{p}\gamma_5\Big]. \tag{2}$$

Here the structure of Dirac matrices is determined by Lorentz in-variance and parity. In this paper, we consider the behavior of this correlation function (2) at the kinematical point where the four momenta of N and N^* are both fixed at p_μ and in the soft pion limit $q \to 0$. In this choice, symmetry structure of the correlation function is best studied, which is our main interest in the present paper.

In general, the correlation function (2) contains contributions from $g_{\pi NN}$, $g_{\pi NN^*}$, and $g_{\pi N^*N^*}$. However, it is possible to show that the information of $g_{\pi NN^*}$ can be

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extracted from the $\not p \gamma_5$ term. Using the phenomenological πNN^* Lagrangian

$$\mathcal{L}_{\pi NN^*} = g_{\pi NN^*} \bar{N}^* \tau^a \pi_a N + (\text{H.C.}), \tag{3}$$

the πNN^* contribution in the $\Pi^\pi(p)$ is given by

$$g_{\pi NN^*} \lambda_N(t) \lambda_{N^*}(s) \left[\frac{p^2 + m_N m_{N^*}}{(p^2 - m_N^2)(p^2 - m_{N^*}^2)} + \frac{\not p (m_N + m_{N^*})}{(p^2 - m_N^2)(p^2 - m_{N^*}^2)} \right] i\gamma_5 . \tag{4}$$

Here we have picked up the term where $J(x;t)$ couples to N and $J(x;s)$ to N^*, corresponding to the physical process $N^* \to N + \pi$. The coupling strengths λ are defined by $\langle 0 | J(x;t) | N \rangle = \lambda_N(t) u_N(x)$ and $\langle 0 | J(x;s) | N^* \rangle = \lambda_{N^*}(s) i\gamma_5 u_{N^*}(x)$. Thus, the πNN^* coupling appears in both γ_5 and $\not p \gamma_5$ terms. On the other hand, using the phenomenological πNN interaction Lagrangian

$$\mathcal{L}_{\pi NN} = g_{\pi NN} \bar{N} i\gamma_5 \tau^i \pi^i N , \tag{5}$$

the πNN contribution is given by

$$g_{\pi NN} \lambda_N(t) \lambda_N(s) \frac{i\gamma_5}{p^2 - m_N^2} . \tag{6}$$

The form of Eq. (6) applies also to the $\pi N^* N^*$ coupling. Therefore, the γ_5 term contains information of $\pi NN, \pi N^* N^*$, and πNN^* couplings, while the $\not p \gamma_5$ term contains the information only of the πNN^* coupling.

In recent reports [10], we have computed the two-point correlation function (2) in the operator product expansion (OPE) and found that the $\not p \gamma_5$ term vanishes up to order dimension eight. We have demonstrated that this result is a consequence of chiral symmetry of the interpolating field $J(x)$. Recently, Cohen and Ji have classified various hadron interpolating fields based on chiral symmetry [11]. Such a classification is useful for the study of model independent properties. We here propose new model independent predictions regarding the meson-baryon couplings based on chiral symmetry.

Now let us briefly review how the $\not p \gamma_5$ term for $g_{\pi NN^*}$ vanishes. First we rewrite the correlation function (2) in terms of the communication relation with the isovector axial charge Q_A^a:

$$\Pi^{\pi^a}(p) = \lim_{q \to 0} \int d^4x\, e^{ipx} \langle 0 | TJ(x) \bar{J}(0) | \pi^a(q) \rangle = -\frac{i}{\sqrt{2} f_\pi} \int d^4x\, e^{ipx} \langle 0 | [Q_A^a, TJ(x)\bar{J}(0)] | 0 \rangle$$
$$= -\frac{i}{2\sqrt{2} f_\pi} \int d^4x\, e^{ipx} \{ \gamma_5 \tau^a, \langle 0 | TJ(x) \bar{J}(0) | 0 \rangle \} . \tag{7}$$

Here we have used the transformation property of the interpolating field J:

$$[Q_A^a, J] = \frac{1}{2} \gamma_5 \tau^a J . \tag{8}$$

Using the Dirac structure of the vacuum to vacuum matrix element of (7)

$$\int d^4x\, e^{ipx} \langle 0 | TJ(x) \bar{J}(0) | 0 \rangle \sim A \not p + B 1 , \tag{9}$$

we find that the $\not p \gamma_5$ term disappears in (7).

We should make one remark here. When we write the phenomenological correlation function (4), only one term for $N^* \to N + \pi$ has been considered, whereas in the theoretical expression (2) another contribution from the reversed process $N + \pi \to N^*$ is also contained. If both the contributions are included, the phenomenological expression for the $\not p \gamma_5$ term is factorized by $\lambda_N(t)\lambda_{N^*}(s) - \lambda_N(s)\lambda_{N^*}(t)$. On the other hand, in the OPE side, the $\not p \gamma_5$ term has the common factor $s - t$ [10]. Thus in both cases the correlation function vanishes trivially when $s = t$. However, in Eq. (7) the $\not p \gamma_5$ term vanishes not by this trivial factor but due to the chiral symmetry of the interpolating field. This is what we emphasize in this paper.

The key of the above proof is that Eq. (8) is satisfied regardless of the choice of t, as the nucleon interpolating field $J(x;t)$ transforms as the fundamental representation of the chiral group SU $(2)_R \times$ SU $(2)_L$ To look at this point in some detail, let us investigate the algebraic structure of the interpolating field. The nucleon field which consists of three quarks belongs to the following irreducible representation of SU $(2)_L \times$ SU $(2)_R$:

$$\left[\left(\frac{1}{2},0\right) + \left(0,\frac{1}{2}\right)\right]^3 = \left[\left(\frac{3}{2},0\right) + \left(0,\frac{3}{2}\right)\right] + 3\left[\left(\frac{1}{2},1\right) + \left(1,\frac{1}{2}\right)\right] + 3\left[\left(\frac{1}{2},\tilde{0}\right) + \left(\tilde{0},\frac{1}{2}\right)\right]$$
$$+ 2\left[\left(\frac{\tilde{1}}{2},0\right) + \left(0,\frac{\tilde{1}}{2}\right)\right], \tag{10}$$

where, according to Ref. [11], tildes imply that a pair of left or right quarks are coupled to the isospin singlet. The relevant terms for the nucleon are then $(\frac{1}{2},\tilde{0}) + (\tilde{0},\frac{1}{2})$ and $(\frac{\tilde{1}}{2},0) + (0,\frac{\tilde{1}}{2})$, which correspond to two independent terms of the interpolating field, $J(x;t=-1)$ and $J(x;t=1)$, respectively. However, these two terms cannot be distinguished within SU $(2)_R \times$ SU $(2)_L$, as they carry the same SU (2) axial charge. This is the underlying reason that the commutation relation (8) holds regardless of the parameter t.

Now me mention the ηNN^* coupling briefly. Since the η meson is an isospin singlet, we investigate the U $(1)_A$ property of the nucleon. The nucleon field $J(x,t)$ transforms under U $(1)_A$ transformations as follows:

$$[Q_A, J(x;t=-1)] = \gamma_5 J(x;t=-1)\,, \qquad (11)$$

$$[Q_A, J(x;t=1)] = 3\gamma_5 J(x;t=1)\,, \qquad (12)$$

where Q_A is the U $(1)_A$ axial charge. A crucial point here is that the transformation rule of J differs for different values of t, and, therefore, the correlation function for η cannot reduce to the anticommutation relation with γ_5 unlike Eq. (7) for the pion. Thus the $g_{\eta NN^*}$ coupling need not vanish.

To summarize briefly, we have shown that the symmetry properties of the interpolating field J lead to the suppression of $g_{\pi NN^*}$, while $g_{\eta NN^*}$ is not subject to the similar constraint. Phenomenologically, these properties seem to be well satisfied by the negative parity nucleon N (1535).

The situation becomes less trivial for the three flavor case of SU $(3)_R \times$ SU $(3)_L$. The reason is that while the SU (3) baryons belong to an octet representation of the diagonal vector group SU $(3)_V$, there behavior under axial transformations is not uniquely determined. But once again, for nonstrange nucleons, we find similar results for meson-baryon couplings.

The decomposition of baryon interpolating fields consisting of three quarks of fundamental representations of SU $(3)_R \times$ SU $(3)L$ is accomplished as

$$[(3,1)+(1,3)]^3 = [(10,1)+(1,10)] + 3\,[(6,3)+(3,6)] + 3\,[(3,\bar{3})+(\bar{3},3)] + 2\,[(8,1)+(1,8)] + [(\tilde{1},1)+(1,\tilde{1})]\,. \qquad (13)$$

Here the multiplets assigned to the spin $\frac{1}{2}$ octet baryons corresponding to Eq. (1) are $(8,1)+(1,8)$ and $(3,\bar{3})+(\bar{3},3)$, as they both transform as an octet representation under SU $(3)_V$ transformations. In contrast, they transform in complex ways under SU (3) axial transformations. Explicitly, denoting the baryon interpolating field which belongs to the SU (3) multiplet (p,q) by $J_a^{(p,q)}(a=1,2,\ldots,8)$, the transformation rule is

$$[Q_A^a, J_b^{(8,1)}] = i f_{abc} J_c^{(8,1)}, \qquad (14)$$

$$[Q_A^a, J_b^{(1,8)}] = -i f_{abc} J_c^{(1,8)}, \qquad (15)$$

$$[Q_A^a, J_b^{(3,\bar{3})}] = d_{abc} J_c^{(3,\bar{3})}, \qquad (16)$$

$$[Q_A^a, J_b^{(\bar{3},3)}] = -d_{abc} J_c^{(\bar{3},3)}, \qquad (17)$$

where f_{abc} and d_{abc} are the structure constants of SU(3).

Consider, for example, the transition $p^* \to p\pi^0$. The nucleon interpolating field is superposition of $(8,1)+(1,8)$ and $(3,\bar{3})+(\bar{3},3)$ with the parameter $\alpha(=\frac{1-t}{1+t})$:

$$J_a(\alpha) = J_a^8 + \alpha J_a^3\,, \qquad (18)$$

$$J_a^8 \equiv J_a^{(8,1)} + J_a^{(1,8)}\,, \qquad (19)$$

$$J_a^3 \equiv J_a^{(\bar{3},3)} + J_a^{(3,\bar{3})}\,. \qquad (20)$$

Thus the correlation function takes the form

$$\Pi_{ab} = \langle J_a(\alpha), \bar{J}_b(\beta)\rangle = \langle J_a^8 + \alpha J_a^3, \bar{J}_b^8 + \beta \bar{J}_b^3\rangle\,. \qquad (21)$$

Since the flavor structure of the proton is $p \sim \lambda^4 + i\lambda^5$, we investigate the commutation relations of the 44, 45, 54, and 55 components and $Q_A^{a=3}$. After some computation, we find that

$$\begin{aligned}[Q_A^3, \Pi_{pp}] &\sim [Q_A^3, \Pi_{4+i5,4+i5}] \\ &= [Q_A^3, \Pi_{44} + i\Pi_{54} - i\Pi_{45} + \Pi_{55}] \\ &= \frac{1}{2}(\{i\gamma_5, \Pi_{pp}^{(88)}\} + \alpha\{i\gamma_5, \Pi_{pp}^{(83)}\} \\ &\quad + \beta\{i\gamma_5, \Pi_{pp}^{(38)}\} + \alpha\beta\{i\gamma_5, \Pi_{pp}^{(33)}\})\,, \end{aligned} \qquad (22)$$

where $\Pi_{pp}^{88} = \langle J_p^8, J_p^8\rangle$. Thus we have found again anticommutators with γ_5 as in Eq. (7). There is no difficulty to see the similar result for the neutron. As in the preceding discussion, this result for the proton and neutron follows completely from the symmetry property of (p,n) in the SU $(2)_R \times$ SU $(2)_L$ subgroup of SU $(3)_R \times$ SU $(3)_L$. Indeed, by identifying (8,1) with $(\frac{\tilde{1}}{2},0)$, and $(3,\bar{3})$ with $(\frac{1}{2},\tilde{0})$, one can verify that both J_N^8 and J_N^3 $(N=p,n)$ carry the same SU (2) axial charge.

For the coupling with η which is now a combination of a singlet η_0 and an octet η_8, we must look at the axial charges Q_A^0 and Q_A^8. It is easy to verify that both axial charges differ for J_N^8 and J_N^3, and that the transformation rules of J_N^8 and J_N^3 cannot be written in a unique way as in Eq. (8). Therefore, once again, we find that the $g_{\eta NN^*}$ coupling need not vanish.

One may wonder whether another isospin doublet (Ξ^0, Ξ^-) shares similar properties for coupling constants. However, axial transformations of J_Ξ under SU $(2)_R \times$ SU $(2)L$ are different from those of J_N and we find that for $a = 1, 2, 3$

$$[Q_A^a, J_\Xi^8] = \frac{1}{2}\gamma_5\tau^a J_\Xi^8 , \tag{23}$$

$$[Q_A^a, J_\Xi^3] = -\frac{1}{2}\gamma_5\tau^a J_\Xi^3 . \tag{24}$$

These equations show that J_Ξ^8 and J_Ξ^3 carry axial charges withopposite signs. Thus for (Ξ^0, Ξ^-), for coupling of, for example, $(\Xi^0)^* \rightarrow \Xi^0 + \pi^0$ need not vanish.

One can generalize the above argument to other SU (2) subgroups of SU (3), as corresponding to the U and V spins. Here the baryons are identified with fundamental representations of the SU (2) subgroups, while mesons with adjoint representations. For the U spin, the two baryon doublets are (Σ^-, Ξ^-) and (p, Σ^+), and the meson triplet is $(K^0, \bar{K}^0, \bar{d}d - \bar{s}s)$. Thus, the coupling constant, for instance, for $(\Sigma^-)^* \rightarrow \Xi^- + K^0$ vanishes, while that of $p^* \rightarrow \Sigma^+ + K^0$ need not. Similarly, for the V spin, the relevant baryons and mesons are (Σ^+, Ξ^0), (n, Σ^-), and $(K^+, K^-, \bar{u}u - \bar{s}s)$. It would be interesting if we could see how the U and V spin symmetries are realized in terms of these meson-baryon coupling constants.

Finally, we comment on the recent work of Kim and Lee [12] from a point of view of the chiral symmetry. For the negative parity nucleon, they have adopted an alternative interpolating field which involves a derivative, whose coupling structure is given by

$$\langle 0|J_{N^*}|N^*\rangle = i\lambda_{N^*}\gamma_5 z_\mu \gamma^\mu u_{N^*} . \tag{25}$$

Here z_μ is an auxiliary spacelike vector which is orthogonal to the four momentum carried by the resonance state. Adopting this interpolating field with derivatives for N^*, the chirality of the negative parity nucleon differs from that of the positive parity nucleon due to the presence of the γ_μ matrices in (25). This means that the axial charges of the positive and negative parity nucleons are no longer the same, and, therefore, our argument for vanishing coupling constants cannot be applied any more here [13].

In summary, we have investigated positive and negative parity baryons (B and B^*) produced by the interpolating field without derivatives. We found that the chiral structure of B and B^* determines the properties of meson-BB^* couplings and that they vanish if B and B^* carry the same axial charge. For the πNN^* coupling the SU (2) triplet axial charges are the same for N and N^*, while for the ηNN^* the singlet axial charges are different for N and N^*. This leads to the suppression of the πNN^* coupling while there is no such suppression for ηNN^*. In the real world, the negative parity nucleon $N(1535)$ seems to satisfy these properties reasonably well. Finally, we have extended this argument to SU (3) baryons and shown that which meson-BB^* couplings vanish. It will be interesting if these predictions can be studied experimentally.

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*Electronic address: jido@th.phys.titech.ac.jp

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