Chiral Symmetry for Positive and Negative Parity Nucleons

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Chiral properties of positive and negative parity nucleons, $N$ and $N^*$, are studied from the viewpoint of chiral symmetry. Two possible ways to assign chiral transformations to the negative parity nucleon are considered. Using linear sigma models based on the two chiral realizations, theoretical as well as phenomenological consequences of the two different assignments are investigated. We find that the nucleon mass in the chiral restored phase is the key quantity to determine the meson-nucleon couplings and the axial charges of nucleons. We also discuss the role of chiral symmetry breaking in the mass splitting of $N$ and $N^*$ in the two sigma models.

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Chiral symmetry and its spontaneous breakdown are very important to understand properties of hadrons at low energy. While the spectrum of observed hadrons does not respect chiral symmetry, the broken chiral symmetry is believed to be restored at finite temperature and/or at finite density. In the chiral restored phase, we expect that hadrons in the same irreducible representation of the chiral group are degenerate. Hadrons in such a multiplet are called chiral partners to each other. So far the role of the chiral symmetry has been extensively worked out for meson properties \cite{DK}. For instance, $(\sigma, \pi)$ and $(\rho, a_1)$ are candidates of chiral partners, as they belong to the $(\frac{1}{2}, \frac{1}{2})$ and $(1,1)$ representations of the chiral $SU(2)_L \times SU(2)_R$ group, respectively. In contrast, chiral symmetry of baryons is less understood. Chiral properties of hadrons are very important in the study of, for instance, QCD phase transitions, since transition properties depend crucially on the particle spectrum before and after the transition.

So far there are not many works investigating negative parity nucleons from the point of view of chiral symmetry \cite{DeTar, Matsu}. DeTar and Kunihiro (DK) studied positive and negative parity nucleons in an extended $SU(2)$ linear sigma model \cite{DeTar}. Their work was motivated by the lattice QCD observations which indicated existence of finite mass nucleons after the chiral symmetry is restored \cite{DK}. A similar observation was also made by Schäfer and Shuryak using an instanton liquid model \cite{Sch}. When chiral symmetry is restored, one would naively expect that nucleons become massless as is indicated in the linear sigma model. However, they showed that in the presence of positive and negative parity nucleons it is possible to construct a theory which allows finite nucleon masses without destroying chiral symmetry. Some of nucleon properties then depend crucially on the way how chiral symmetry is implemented.

Recently we have studied properties of negative parity baryons using the QCD sum rule approach \cite{DDK}. We have calculated the masses of positive and negative parity baryons ($B_+$ and $B_-$) in the flavor octet and singlet representations. We have found that the quark condensates, which break chiral symmetry, induce the mass splitting between $B_+$ and $B_-$. When the chiral order parameter $\langle \bar{q}q \rangle$ vanishes, we have found that $B_-$ is degenerate with $B_+$ and also that they tend to become massless. This result may be contrasted with the observation made by DeTar and Kunihiro. We have also investigated the $\pi NN^*$ coupling in the same framework of the QCD sum rule \cite{DDK}. (Throughout this paper $N^*$ denotes a negative parity nucleon, e.g. $N(1535)$.) There it is found that the $\pi NN^*$ coupling vanishes in the chiral and soft-pion limit. On the other hand, as opposed to our finding, Kim and Lee have obtained a non-vanishing $\pi NN^*$ coupling when an alternative interpolating field that contains a derivative is used for $N^*$ \cite{KimLee}.

The purpose of the present letter is to point out the importance of chiral symmetry for properties of positive and negative parity nucleons. It turns out that when we treat two kinds of nucleons, there are two ways to assign chiral transformations of baryons. As a consequence we can construct different chiral effective models based on the two chiral realizations, and then we clarify why different results were obtained from various approaches discussed above. We study masses, $\pi NN^*$ couplings and axial charges of $N$ and $N^*$. We also discuss the role of the chiral symmetry breaking in the $N-N^*$ mass splitting. To accomplish our purpose, we first consider chiral transformations of $N$ and $N^*$, and then by using the linear sigma model, we discuss the behaviors of $N$ and $N^*$ under the chiral phase transition.

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Let us start with considering the chiral transformation for one nucleon $N$. To be definite we will consider the $SU(2)_R \times SU(2)_L$ chiral group. We assume that the nucleons belong to linear representations of the chiral group. This may not be the case when chiral symmetry is spontaneously broken, where nonlinear representations are also possible. In that case, however, it is difficult to study the transition properties of the chiral symmetry. Therefore, we will adopt linear representations for a consistent study of chiral restoration properties.

The chiral transformation for $N$ is defined by

$$N_R \rightarrow RN_R, \quad N_L \rightarrow LN_L,$$

where $R$ ($L$) is an element of $SU(2)_R$ ($SU(2)_L$), and $N_R$ ($N_L$) are the right (left) component of the Dirac spinors, satisfying $\gamma_5 N_R = N_R$ and $\gamma_5 N_L = -N_L$. Eq. (1) is no more than definition; the transformation for the “right” (“left”) handed nucleon is just called the “right” (“left”) transformation.

Chiral symmetry does not allow a mass term for $N$ in the Lagrangian, so that $N$ is a massless particle in the Wigner phase (the chiral restored phase) and the nucleon mass is generated by spontaneous chiral symmetry breaking. To find the chiral partner of $N$, we consider the commutation relation of the generators of the axial transformation $X^a (a = 1, 2, 3)$ and $N$ in the Wigner phase:

$$[X^a, N] = -i\gamma_5 \frac{\tau^a}{2} N,$$

where $\tau^a$ is the isospin matrix. This commutation relation follows from the chiral transformation (1). The commutation relation (2) implies that the chiral partner of $N = N_R + N_L$ is $\gamma_5 N = N_R - N_L$ and that no additional particle is necessary to complete the chiral multiplet $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$. It should be noted that because $N$ is massless in the Wigner phase, we may regard $N_R$ and $N_L$ as independent particles. In other words a massive fermion in NG phase (the chiral broken phase), which has four components, splits into two massless fermions with two components each in the Wigner phase and they form the pair of the chiral multiplet.

Next, we consider additional nucleon $N^*$ with negative parity. First we introduce two nucleon fields $N_1$ and $N_2$ each of which belongs to the chiral multiplet $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$. The physical nucleons $N$ and $N^*$ are linear combinations of $N_1$ and $N_2$, when the Lagrangian has the mixing terms. As anticipated, there are two possible assignments of chiral transformations. In the first scheme, which we call the “naive assignment”, both $N_1$ and $N_2$ transform in the same way. In the second scheme, which we call the “mirror assignment”, the second nucleon transforms in the reversed way to the first nucleon.

In the naive assignment the chiral transformation for $N_1$ and $N_2$ is defined by

$$N_{1R} \rightarrow RN_{1R}, \quad N_{1L} \rightarrow LN_{1L},$$

$$N_{2R} \rightarrow RN_{2R}, \quad N_{2L} \rightarrow LN_{2L}.$$  

Chiral symmetry requires again that these two nucleons must be massless in the Wigner phase. In order to find their chiral partners, we consider the commutation relations which follow from (3) and (4):

$$[X^a, N_1] = -\frac{1}{2} i\gamma_5 \tau^a N_1, \quad [X^a, N_2] = -\frac{1}{2} i\gamma_5 \tau^a N_2.$$  

These relations show that $N_1$ and $N_2$ belong to the multiplet $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$ separately. Therefore, as in the previous case, the chiral partner of $N_1$ is $\gamma_5 N_1$, and that of $N_2$ is $\gamma_5 N_2$. This situation does not change even if many nucleons with the naive transformation are considered.

We note that the commutation relations (3) imply that $g_A = 1$ for both $N$ and $N^*$ in the linear sigma model unless we introduce derivative couplings of $\sigma$ and $\pi$ with the nucleon. It is important to note that the sign of the axial charge $g_A$ of $N^*$ is the same as the one of $N$. As we will see later, in the mirror case the axial charge of $N^*$ has the opposite sign to $N$. This is notable difference between the naive model and the mirror model.

In the mirror assignment, denoting the two nucleon fields by $\psi_1$ and $\psi_2$, the transformation rule is defined as

$$\psi_{1R} \rightarrow R\psi_{1R}, \quad \psi_{1L} \rightarrow L\psi_{1L},$$

$$\psi_{2R} \rightarrow L\psi_{2R}, \quad \psi_{2L} \rightarrow R\psi_{2L}.$$  

The right (left) component of $\psi_2$ transforms as the left (right) component of $\psi_1$. The reason that this assignment is possible is that the left- and right-handedness of the fermion, which is determined by the eigenvalue of $\gamma_5$, is independent of that of chiral symmetry, although we use the same terminology. The chirality of fermion specifies representations of the Lorentz group for the fermion while the chirality of chiral symmetry is associated with the internal chiral symmetry $SU(2)_R \times SU(2)_L$. In this case, we are allowed to introduce a chirally invariant mass term.
Therefore, the nucleons can have a finite mass \( m_0 \) when the chiral symmetry is restored.

Using eqs. (8) and (7), we find the following commutation relations

\[
[X^a, \psi_1] = -i\gamma_5 \frac{\tau^a}{2} \psi_1, \\
[X^a, \psi_2] = +i\gamma_5 \frac{\tau^a}{2} \psi_2.
\]

(9)

Note that the sign on the r.h.s. for \( \psi_2 \) is opposite to that for \( \psi_1 \). In order to obtain the physical nucleons \( \psi_+ \) and \( \psi_- \), we have to diagonalize the mass term (8) because it is off-diagonal in the basis \((\psi_1, \psi_2)\). In the diagonalized basis, the commutation relations are given by

\[
[X^a, \psi_+] = -\frac{\sigma_0}{2} \psi_-, \\
[X^a, \psi_-] = +\frac{\sigma_0}{2} \psi_+,
\]

(10)

where \( \psi_+ = \frac{1}{\sqrt{2}}(\psi_1 + \gamma_5 \psi_2) \) and \( \psi_- = \frac{1}{\sqrt{2}}(\gamma_5 \psi_1 - \psi_2) \) in the Wigner phase. From the commutation relation (10) we see that \( \psi_+ \) and \( \psi_- \) are transformed into each other under the chiral transformation, and therefore, \( \psi_+ \) and \( \psi_- \) belong to the same multiplet of \( SU(2)_R \times SU(2)_L \). In this way, \( \psi_+ \) and \( \psi_- \) are considered to be chiral partners of each other.

More explicitly, in the group theoretical language, \( \psi_{1R} \) and \( \psi_{2L} \) belong to \((\frac{1}{2}, 0)\) and \( \psi_{1L} \) and \( \psi_{2R} \) belong to \((0, \frac{1}{2})\). Because the nucleons have masses in the Wigner phase, we need four components to represent each nucleon. Thus it is appropriate to introduce

\[
\Psi_r = \frac{1}{\sqrt{2}}(\psi_{1R} \pm \psi_{2L})
\]

(11)

\[
\Psi_l = \frac{1}{\sqrt{2}}(\psi_{2R} \pm \psi_{1L})
\]

(12)

which have four independent components each. We see that \( \Psi_r \) belongs to \((\frac{1}{2}, 0)\) and \( \Psi_l \) belongs to \((0, \frac{1}{2})\) and that the parity eigenstates \( \psi_+ \) and \( \psi_- \) are given by linear combinations of \( \Psi_r \) and \( \Psi_l \). Thus \( \psi_+ \) and \( \psi_- \) belong to the same multiplet \((\frac{1}{2}, 0) \oplus (0, \frac{1}{2})\) and the chiral partner of \( \psi_+ \) is \( \psi_- \).

In order to see the above argument more concretely, we now construct linear sigma models for the two kinds of nucleons which satisfy the above chiral symmetries.

First we consider the case of the naive assignment. Considering the transformation rule for the meson field \( M \equiv \sigma + i \vec{\pi} \cdot \vec{\pi} \rightarrow LMR \), we can easily write down a renomalizable chiral invariant Lagrangian:

\[
\mathcal{L}_{\text{naive}} = \tilde{N}_1 \partial \sigma N_1 + \tilde{N}_2 \partial \sigma N_2 + a \tilde{N}_1(\sigma + i \gamma_5 \vec{\pi} \cdot \vec{\pi}) N_1 + b \tilde{N}_2(\sigma + i \gamma_5 \vec{\pi} \cdot \vec{\pi}) N_2 \\
+ c \{ \tilde{N}_2(\gamma_5 \sigma + i \vec{\pi} \cdot \vec{\pi}) N_1 - \tilde{N}_1(\gamma_5 \sigma + i \vec{\pi} \cdot \vec{\pi}) N_2 \} + \mathcal{L}_M,
\]

(13)

where \( a, b \) and \( c \) are coupling constants. Here \( \mathcal{L}_M \) is a chiral invariant meson Lagrangian which is not important in the following discussion. The fifth term in this Lagrangian gives a mixing between \( N_1 \) and \( N_2 \). The chiral symmetry breaks down spontaneously with a finite vacuum expectation value of the sigma meson, \( \sigma_0 \equiv \langle 0 | \sigma | 0 \rangle \). To obtain physical nucleons \( N_+ \) and \( N_- \) in the NG phase, we have to diagonalize the mass matrix \( M \) which is given by

\[
M \sim \sigma_0 \begin{pmatrix}
\frac{a}{\gamma_5 c} & -\gamma_5 c \\
\gamma_5 c & \frac{b}{\gamma_5 c}
\end{pmatrix}
\]

(14)

This matrix can be diagonalized by the physical nucleon fields

\[
\begin{pmatrix}
N_+ \\
N_-
\end{pmatrix} = \frac{1}{\sqrt{2} \cosh \delta} \begin{pmatrix}
e^{\delta/2} & -\gamma_5 e^{-\delta/2} \\
\gamma_5 e^{\delta/2} & -e^{\delta/2}
\end{pmatrix} \begin{pmatrix}
N_1 \\
N_2
\end{pmatrix}
\]

(15)

where the mixing angle \( \delta \) is defined by \( \sinh \delta = -(a + b)/2c \). In this basis the masses of \( N_+ \) and \( N_- \) are given by

\[
m_{\pm} = \frac{\sigma_0}{2} \left( \sqrt{(a + b)^2 + 4c^2} \mp (a - b) \right)
\]

(16)

We present a schematic plot of \( m_{\pm} \) as functions of \( \sigma_0 \) in Fig. We In the Wigner phase, i.e. when \( \sigma_0 \rightarrow 0 \), both \( N_+ \) and \( N_- \) become massless and get degenerate. However, this degeneracy is trivial rather than due to chiral symmetry, because all the nucleons with the naive assignment are massless in the Wigner phase. Similarly, the mass difference
of $N_+$ and $N_-$ is caused by the choice of the coupling parameters, $a$ and $b$ and therefore it is independent of chiral symmetry.

It should be noted that the physical nucleon fields $N_+$ and $N_-$ decouple from each other after the diagonalization, when we truncate the meson-nucleon coupling Lagrangian at the non-derivative Yukawa term, because the coupling term of $\mathcal{L}_\text{DK}$ is factored out by the mass matrix $M$. This implies that the off-diagonal Yukawa coupling $g_{\pi NN}$ vanishes in the soft pion limit, when all the derivative couplings are neglected. This result is qualitatively consistent with the observed $g_{\pi NN}(1535) \sim 1$ which is strongly suppressed in comparison with $g_{\pi NN} \sim 13$. In fact, up to the non-derivative Yukawa coupling, this sigma model is reduced to the sum of two independent sigma model even without the $\sigma$ condensation, since the mixing angle $\delta$ is independent of $\sigma_0$. Namely the parameter $c$, which appears in the off-diagonal Yukawa coupling, is a superficial parameter.

Next we turn to the discussion of the mirror model. In this model the negative parity nucleon is assumed to follow

\[ \mathcal{L}_{\text{DK}} = \bar{\psi}_1 i\partial \psi_1 + \bar{\psi}_2 i\partial \psi_2 + m_0 (\bar{\psi}_2 \gamma_5 \psi_1 - \bar{\psi}_1 \gamma_5 \psi_2) + a \psi_1 (\sigma + i\gamma_5 \vec{\tau} \cdot \vec{\pi}) \psi_1 + b \psi_2 (\sigma - i\gamma_5 \vec{\tau} \cdot \vec{\pi}) \psi_2 + \mathcal{L}_M, \]

where $a = g_2 - g_1$ and $b = -g_1 - g_2$ using the parameters in Refs. $[4][4]$. This Lagrangian was first proposed and studied by DeTar and Kunihiro $[4]$. In the basis $|\psi_1 \rangle$ the masses of $\psi_+$ and $\psi_-$ are given by

\[ m_\pm = \frac{1}{2} \sqrt{(a+b)^2 \sigma_0^2 + 4m_0^2 \mp (a-b)\sigma_0}, \]

A schematic plot of $m_\pm$ as functions of $\sigma_0$ is presented in Fig. 1. In the Wigner phase these nucleons are degenerate with a finite mass $m_0$. This shows the nucleon masses are generated by $m_0$, which is quite different from the mechanism of mass generation in the naive model, where the chiral symmetry breaking creates the nucleon masses. We see that the mass splitting between $\psi_+$ and $\psi_-$ is caused by the spontaneously chiral symmetry breaking. In this sense $m_0$ is the most important parameter in the mirror model. The case $m_0 = 0$ is special, because we can not distinguish the naive model and the mirror model.

We note two more differences in the mirror model from the naive model. First, meson couplings between $\psi_+$ and $\psi_-$ no longer vanish unlike the naive case and can remain finite, because the coupling matrix differs from the mass matrix and need not be diagonalized in the basis $|\psi_+, \psi_-\rangle$ that diagonalizes the mass matrix. Second, the commutation relations between the axial charges $Q_5^a$ and the nucleon fields are quite different,

\[ |Q_5^a, \psi_+\rangle = \frac{\tau^a}{2} (\tanh \delta \gamma_5 \psi_+ + \frac{1}{\cosh \delta} \psi_-) \]

\[ |Q_5^a, \psi_-\rangle = \frac{\tau^a}{2} (-\tanh \delta \gamma_5 \psi_- + \frac{1}{\cosh \delta} \psi_+) \]

They imply that the axial charges are now given in the form of a $2 \times 2$ matrix and that the sign of $g_A^{\psi_+ \psi_-}$ is opposite to $g_A^{\psi_- \psi_+}$. It would be of great interest to see the relative sign of the axial charges experimentally, as it provides the key information on the chiral structure of negative parity nucleon. We summarize a comparison between the naive and mirror model in Table 1.

Finally, we comment on the QCD sum rule analysis on the negative parity nucleon. It turns out that the chiral assignment in our QCD sum rule $[8][8]$ corresponds to the naive assignment. In this analysis we have introduced $N$ and $N^*$ as

\[ \langle 0 | J(x) | N \rangle = \lambda_N u_N(x), \]

\[ \langle 0 | J(x) | N^* \rangle = i \gamma_5 \lambda_{N^*} u_{N^*}(x). \]
Then the same chiral transformation for \( N^* \) as for \( N \) follows. This is the basis of our previous QCD sum rule results, i.e., that \( N \) and \( N^* \) tend to become massless as \( \langle \bar{q}q \rangle \to 0 \) and that the \( \pi NN^* \) coupling vanishes in the soft and chiral limit.

On the other hand, Kim and Lee found quite different results from ours [9]. While they have used the same type of interpolating field for \( N \) as our choice, for \( N^* \) they have adopted an alternative interpolating field \( \eta_{N^*} \) that contains a derivative [11]. Although the \( \eta_{N^*} \) itself transforms in the same way as \( N \), the chiral structure of \( N^* \) is changed by the coupling of \( N^* \) to \( \eta_{N^*} \):

\[
\langle 0 | \eta_{N^*} | N^* \rangle = i\lambda_{N^*} \gamma_5 z_\mu \gamma_\mu u_{N^*},
\]

where \( z_\mu \) is an auxiliary space-like vector which is orthogonal to the four momentum carried by the resonance state.

The \( \gamma_\mu \) matrix on the r.h.s. of (24) changes the chirality and thus makes \( N^* \) being a mirror of \( N \). This is the reason why they have obtained a finite \( \pi NN^* \) coupling.

We need further investigations on chiral properties of baryons, because we know very little which of the two assignments is realized in the physical nucleons. As far as the sum rule analysis are compared, there is no strong preference of one of the two assignments. DeTar and Kunihiro reproduced the masses of \( N \) and \( N(1535) \) and the observed \( g_{\pi NN(1535)} \) by choosing \( m_0 = 270 \) MeV [4]. The DK model and our extended model to \( SU(3) \) [10] suggest that the suppression of the \( \pi NN^* \) coupling is caused by the smallness of \( m_0 \), which is the key parameter characterizing this model. For small \( m_0 \), the mirror assignment may give similar predictions for the masses of \( N \) and \( N^* \) as well as the \( \pi NN^* \) coupling strengths and therefore can hardly be distinguished phenomenologically. However, a notable difference between the two choices is the sign of the axial charge of \( N^* \), \( g_{N^*N^*A} \). It would be extremely interesting if the axial charge of \( N^* \) is observed experimentally.

In summary, we have investigated the properties of the negative parity nucleon \( N^* \) from the viewpoint of chiral symmetry. We have two ways of assignments of chiral transformation for \( N^* \), so that we obtain two linear sigma models based on them. We have observed several qualitative differences of the properties of \( N \) and \( N^* \) between the two, which are summarized in Table I. The origin of the differences is the chiral structures of the nucleons. In the naive case, \( N^* \) has nothing to do with \( N \) in the sense that the \( N^* \) belongs to different multiplets from the one of \( N \), while in the mirror case \( N^* \) belongs to the same multiplet as that of \( N \), so that \( N^* \) can be interpreted as the chiral partner of \( N \). If in the real world the mirror case is realized, we need to identify the chiral partner of \( N(939) \) with the negative parity nucleon such as \( N(1535) \) or \( N(1650) \).

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### Table I.

| Definition                        | Naive Assignment | Mirror Assignment |
|-----------------------------------|------------------|-------------------|
| \( \pi NN^* \) coupling          | \( N_{2R} \to RN_{2R} \), \( N_{2L} \to LN_{2L} \) | \( \psi_{2R} \to L\psi_{2R} \), \( \psi_{2L} \to R\psi_{2L} \) |
| Chiral partner                    | \( N_+ \leftrightarrow \gamma_5 N_+ \), \( N_- \leftrightarrow \gamma_5 N_- \) | \( (a + b)/cosh \delta \) |
| \( g_A \) \( g_{NN^*} \)         | Positive mass generation | \( \psi_+ \leftrightarrow \psi_- \) |
| Role of \( \sigma_0 \)            | Positive mass generation | Negative mass splitting |
FIG. 1. A schematic plot of $\sigma_0$ dependences of $N$ and $N^*$ masses in the naive model (the solid line) and the mirror model (the dashed line).

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