Physics Letters B 611 (2005) 87–92
PHYSICS LETTERS B
www.elsevier.com/locate/physletb

Chiral perturbation theory for pentaquark baryons and its applications

P. Ko [a], Jungil Lee [b], Taekoon Lee [c], Jae-Hyeon Park [a]

[a] Department of Physics, KAIST, Daejon 305-701, South Korea
[b] Department of Physics, Korea University, Seoul 136-701, South Korea
[c] Department of Physics, Seoul National University, Seoul 151-742, South Korea

Received 13 December 2004; accepted 14 February 2005

Editor: G.F. Giudice

Abstract

We construct a chiral Lagrangian for pentaquark baryons assuming that the recently found $\Theta^+(1540)$ state belongs to an antidecuplet of SU(3) flavor symmetry with $J^P = \frac{1}{2}^{\pm}$. We derive the Gell-Mann–Okubo formulae for the antidecuplet baryon masses, and a possible mixing between the antidecuplet and the pentaquark octet. Then we calculate the cross sections for $\pi^- p \to K^- \Theta^+$ and $\gamma n \to K^- \Theta^+$ using our chiral Lagrangian. The resulting amplitudes respect the underlying chiral symmetry of QCD correctly. We also describe how to include the light vector mesons in the chiral Lagrangian.

PACS: 12.39.Fe; 14.20.-c; 13.75.-n; 13.60.-r

1. Introduction

Recently, five independent experiments reported observations of a new baryonic state $\Theta^+(1540)$ with a very narrow width < 5 MeV [1–6], which is likely to be a pentaquark state $(uudd\bar{s})$ [7]. Arguments based on quark models suggest that this state is a member of SU(3) antidecuplet with spin $J = \frac{1}{2}$ or $\frac{3}{2}$. The hadro/photo-production cross section would depend on the spin J and parity P of the Θ^+, and it is important to have reliable predictions for these cross sections. The most proper way to address these issues will be chiral perturbation theory.

In this Letter, we construct a chiral Lagrangian for pentaquark baryons assuming they are SU(3) antidecuplet with $J = \frac{1}{2}$ and $P = +1$ or -1. (The case for $J = \frac{3}{2}$ can be discussed in a similar manner, except that antidecuplets are described by Rarita–Schwinger fields.) Then we calculate the mass spectra of antide-

E-mail addresses: pko@muon.kaist.ac.kr (P. Ko), jungil@hep.anl.gov (J. Lee), tlee@phya.snu.ac.kr (T. Lee), jhpark@muon.kaist.ac.kr (J.-H. Park).

0370-2693
doi:10.1016/j.physletb.2005.02.031

cuplets, their possible mixings with pentaquark octets, the decay rates of antidecuplets, and cross sections for $\pi^- p \to K^- \Theta^+$ and $\gamma n \to K^- \Theta^+$. Finally we describe how to include light vector mesons in our framework, and how the low energy theorem is recovered in the soft pion limit.

2. Chiral Lagrangian for a pentaquark baryon decuplet

Let us denote the Goldstone boson field by pion octet π, baryon octet including nucleons by B, and antidecuplet including Θ^+ by $\mathcal{P}$. Under chiral $SU(3)_L \times SU(3)_R$ [8], the Goldstone boson field $\Sigma \equiv \exp(2i\pi/f)$, where $f \approx 93$ MeV is the pion decay constant, transforms as

$$\Sigma(x) \to L\Sigma(x)R^\dagger.$$

It is convenient to define another field $\xi(x)$ by $\Sigma(x) \equiv \xi^2(x)$, which transforms as

$$\xi(x) \to L\xi(x)U^\dagger(x) = U(x)\xi(x)R^\dagger.$$

The 3×3 matrix field $U(x)$ depends on Goldstone fields $\pi(x)$ as well as the SU(3) transformation matrices L and R. It is convenient to define two vector fields with following properties under chiral transformations:

$$\begin{aligned}
V_\mu &= \frac{1}{2}\left(\xi^\dagger\partial_\mu\xi + \xi\partial_\mu\xi^\dagger\right),\\
V_\mu &\to UV_\mu U^\dagger + U\partial_\mu U^\dagger,\\
A_\mu &= \frac{i}{2}\left(\xi^\dagger\partial_\mu\xi - \xi\partial_\mu\xi^\dagger\right),\\
A_\mu &\to UA_\mu U^\dagger.
\end{aligned} \tag{1}$$

Note that V_μ transforms like a gauge field. The transformation of the baryon octet and pentaquark antidecuplet $\mathcal{P}$ including Θ^+ ($I = 0$) can be chosen as

$$\begin{aligned}
B^i{}_j &\to U^i{}_a B^a{}_b U^{\dagger b}{}_j,\\
\mathcal{P}_{ijk} &\to P_{abc} U^{\dagger a}{}_i U^{\dagger b}{}_j U^{\dagger c}{}_k,
\end{aligned}$$

where all the indices are for SU(3) flavor. The pentaquark baryons are related to $\mathcal{P}_{abc} = \mathcal{P}_{(abc)}$ by, for example, $\mathcal{P}_{333} = \Theta^+$, $\mathcal{P}_{133} = \frac{1}{\sqrt{3}}\tilde{N}^0$, $\mathcal{P}_{113} = \frac{1}{\sqrt{3}}\tilde{\Sigma}^-$, and $\mathcal{P}_{112} = \frac{1}{\sqrt{3}}\Xi^-_{3/2}$. Then, one can define a covariant derivative $\mathcal{D}_\mu$, which transforms as

$$\mathcal{D}_\mu B \to U\mathcal{D}_\mu BU^\dagger, \quad \text{by}$$
$$\mathcal{D}_\mu B = \partial_\mu B + [V_\mu, B].$$

Chiral symmetry is explicitly broken by nonvanishing current-quark masses and electromagnetic interactions. The former can be included by regarding the quark-mass matrix $m = \mathrm{diag}(m_u, m_d, m_s)$ as a spurion with transformation property $m \to LmR^\dagger = RmL^\dagger$. It is more convenient to use $\xi m\xi + \xi^\dagger m\xi^\dagger$, which transforms as an SU(3) octet. Electromagnetic interactions can be included by introducing photon field $\mathcal{A}_\mu$ and its field strength tensor $F_{\mu\nu} = \partial_\mu\mathcal{A}_\nu - \partial_\nu\mathcal{A}_\mu$:

$$\partial_\mu\Sigma \to \mathcal{D}_\mu\Sigma \equiv \partial_\mu\Sigma + ie\mathcal{A}_\mu[Q, \Sigma], \tag{2a}$$

$$V_\mu \to V_\mu + \frac{ie}{2}\mathcal{A}_\mu\left(\xi^\dagger Q\xi + \xi Q\xi^\dagger\right), \tag{2b}$$

$$A_\mu \to A_\mu - \frac{e}{2}\mathcal{A}_\mu\left(\xi^\dagger Q\xi - \xi Q\xi^\dagger\right), \tag{2c}$$

where $Q \equiv \mathrm{diag}(2/3, -1/3, -1/3)$ is the electric-charge matrix for light quarks ($q = u, d, s$).

Now it is straightforward to construct a chiral Lagrangian with lowest order in derivative expansion. The parity and charge-conjugation symmetric chiral Lagrangian is given by

$$\mathcal{L} = \mathcal{L}_\Sigma + \mathcal{L}_B + \mathcal{L}_\mathcal{P}, \tag{3}$$

where

$$\mathcal{L}_\Sigma = \frac{f_\pi^2}{4}\mathrm{Tr}\left[\mathcal{D}_\mu\Sigma^\dagger\mathcal{D}^\mu\Sigma - 2\mu m\left(\Sigma + \Sigma^\dagger\right)\right], \tag{4a}$$

$$\begin{aligned}
\mathcal{L}_B = {}&\mathrm{Tr}\,\bar{B}(i\not{D} - m_B)B + D\,\mathrm{Tr}\,\bar{B}\gamma_5\{\not{A}, B\}\\
&+ F\,\mathrm{Tr}\,\bar{B}\gamma_5[\not{A}, B],
\end{aligned} \tag{4b}$$

$$\begin{aligned}
\mathcal{L}_\mathcal{P} = {}&\bar{\mathcal{P}}(i\not{D} - m_\mathcal{P})\mathcal{P} + \mathcal{C}_{\mathcal{P}N}(\bar{\mathcal{P}}\Gamma_P\not{A}B + \bar{B}\Gamma_P\not{A}\mathcal{P})\\
&+ \mathcal{H}_{\mathcal{P}N}\bar{\mathcal{P}}\gamma_5\not{A}\mathcal{P},
\end{aligned} \tag{4c}$$

where P is the parity of Θ^+, $\Gamma_+ = \gamma_5$, and $\Gamma_- = 1$, and $m_\mathcal{P}$ is the average of the pentaquark decuplet mass.

The Gell-Mann–Okubo formulae for pentaquark baryons will be obtained from

$$\mathcal{L}_m = \alpha_m\bar{\mathcal{P}}\left(\xi m\xi + \xi^\dagger m\xi^\dagger\right)\mathcal{P}. \tag{5}$$

Expanding this, we get the mass splittings $\Delta m_i \equiv m_i - m_\mathcal{P}$ within the antidecuplet:

$$\Delta m_\Theta = 2\alpha_m m_s, \tag{6a}$$

$$\Delta m_{\tilde{N}} = \alpha_m (2\hat{m} + 4m_s)/3, \tag{6b}$$

$$\Delta m_{\tilde{\Sigma}} = \alpha_m (4\hat{m} + 2m_s)/3, \tag{6c}$$

$$\Delta m_{\Xi_{3/2}} = 2\alpha_m \hat{m}, \tag{6d}$$

where $\hat{m} = m_u = m_d$ ignoring small isospin-breaking effects. If the newly observed state at a mass 1862 ± 2 MeV is identified as $\Xi_{3/2}$, we find

$$m_{\tilde{N}} = 1647 \text{ MeV}, \qquad m_{\tilde{\Sigma}} = 1755 \text{ MeV}. \tag{7}$$

Recently Jaffe and Wilczek suggested there could be an ideal mixing between pentaquark antidecuplet $\mathcal{P}_{abc}$ and pentaquark octet $\mathcal{O}^a{}_b$ with the same parities [9]. This idea has been generalized further by other groups [10,11]. In chiral Lagrangian approach, such a general mixing arises from

$$\beta_m \big[\bar{\mathcal{P}}\big(\xi m \xi + \xi^\dagger m \xi^\dagger\big)\mathcal{O} + \bar{\mathcal{O}}\big(\xi m \xi + \xi^\dagger m \xi^\dagger\big)\mathcal{P}\big]. \tag{8}$$

Expanding this leads to

$$\mathcal{L} = \mathcal{B}_m \big[\bar{p}\tilde{N}^+ - \bar{n}\tilde{N}^0 + \bar{\Sigma}^0\tilde{\Sigma}^0 - \bar{\Sigma}^-\tilde{\Sigma}^- + \bar{\Sigma}^+\tilde{\Sigma}^+ + \text{H.c.}\big], \tag{9}$$

where $\mathcal{B}_m = 2\beta_m(m_s - \hat{m})/\sqrt{3}$ and we borrowed baryon-octet notation for pentaquark octet states in Eq. (9). Note that the relative sign between $\bar{n}\tilde{N}^0$ and $\bar{p}\tilde{N}^+$ (and also $\bar{\Sigma}^0\tilde{\Sigma}^0$) is negative, unlike the case in Ref. [10]. This is due to the $\overline{\mathbf{10}}$ nature of the $\mathcal{P}$. One could write down the same mixing between pentaquark antidecuplet $\mathcal{P}_{abc}$ and the ordinary baryon octet B, but such terms will be highly suppressed compared to the above term, since it is a mixing between qqq and $qqqq\bar{q}$.

Finally the baryon decuplet can only couple to pentaquark octet $\mathcal{O}$, but not to pentaquark antidecuplet $\mathcal{P}$, since $\mathbf{10} \otimes \mathbf{8} \otimes \mathbf{10}$ does not contain SU(3) singlet. This implies that $N(1440)$ or $N(1710)$ cannot be pure pentaquark antidecuplets, because they have substantial branching ratios into $\Delta\pi$ final states. They could be mixed states of pentaquark octet and pentaquark antidecuplet, and their productions and decays will be more complicated than pure antidecuplet case. Since the current data on baryon sectors are not enough to study such mixings in details, we do not pursue the mixing further in the following.

Parameters in the above Lagrangian are taken to have the following numerical values: $m_B \approx 940$ MeV is the nucleon mass, $D \approx -0.81$ and $F \approx -0.47$ at tree level, and we assume $\hat{m} = 0$ and $m_\eta^2 = (4/3)m_K^2$.

The coupling $\mathcal{C}_{\mathcal{P}N}$ is determined from the decay width Γ_Θ of the Θ^+, which is dominated by K^+n and K^0p modes as $\Gamma_\Theta/2 = \Gamma_{\Theta^+\to K^+n} = \Gamma_{\Theta^+\to K^0p}$:

$$\Gamma_\Theta = \frac{\mathcal{C}_{\mathcal{P}N}^2 |\mathbf{p}^*|}{8\pi f^2 m_\Theta^2}(m_\Theta \pm m_B)^2\big[(m_\Theta \mp m_B)^2 - m_K^2\big], \tag{10}$$

where $\mathbf{p}^*$ is the kaon momentum in the Θ^+ rest frame and the signs are for $P(\Theta^+) = \pm 1$, respectively. Then the $\mathcal{C}_{\mathcal{P}N}$ is determined as

$$\mathcal{C}_{\mathcal{P}N}^2(P = +, -) = (2.7, 0.90) \times \Gamma_\Theta/\text{GeV}.$$

Cahn and Trilling [12] argues that $\Gamma_\Theta = (0.9 \pm 0.3)$ MeV using the DIANA results [2]. We present our results proportional to the $\mathcal{C}_{\mathcal{P}N}^2$ rescaled by the factor $1 \text{ MeV}/\Gamma_\Theta$. Such a small $\mathcal{C}_{\mathcal{P}N}$ can be understood as following: this coupling is related to the matrix element of hadronic axial vector current operator (with zero baryon number) between a pentaquark baryon and an ordinary baryon. Since two states have different number of valence quarks, this matrix element should be highly suppressed compared to the ordinary axial vector coupling $(D + F)$ or $\mathcal{H}_{\mathcal{P}N}$.

The coupling $\mathcal{H}_{\mathcal{P}N}$ is also unknown, and determines transition rates between pentaquark antidecuplets with pion or kaon emission. Unfortunately, such decays are all kinematically forbidden, and cannot be used to fix $\mathcal{H}_{\mathcal{P}N}$. However, we expect that $\mathcal{H}_{\mathcal{P}N} = O(1)$, without any suppression as in $\mathcal{C}_{\mathcal{P}N}$. With this remark in mind, we will assume $\mathcal{H}_{\mathcal{P}N}$ can vary between -4 and 4 in the numerical analysis to be as general as possible.

3. $\pi^- p \to K^- \Theta^+$

Let us first consider $\pi^- p \to K^- \Theta^+$ as applications of our chiral Lagrangian for pentaquark baryons. In Fig. 1(a) and (b), we show the relevant Feynman diagrams. Note that only Fig. 1(a) was considered in the literature. However, there is an s-channel $\tilde{N}^0(1647)$ exchange diagram Fig. 1(b) in our chiral Lagrangian, since Θ^+ is not an SU(3) singlet, but belongs to the antidecuplet. Therefore one has to keep both Fig. 1(a) and (b) in order to get an amplitude with correct SU(3) flavor symmetry.

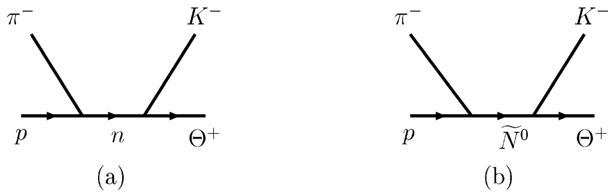

Fig. 1. Feynman diagrams for $\pi^- p \to K^- \Theta^+$.

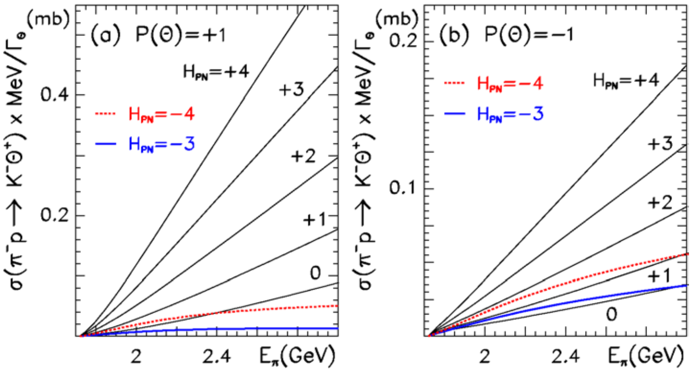

Fig. 2. Cross sections for $\pi^- p \to K^- \Theta^+$ for selective values of $\mathcal{H}_{\mathcal{P}N}$ between -4 and $+4$, in the proton rest frame: (a) $P(\Theta) = +1$ and (b) $P(\Theta) = +1$.

The amplitude for $\pi^- p \to K^- \Theta^+$ is given by

$$\mathcal{M} = \frac{\mathcal{C}_{\mathcal{P}N}}{2f^2} \bar{u}_{\Theta^+} \left(\Gamma_P \not{p}_{K^-} \frac{D+F}{\not{p}_{\pi^-} + \not{p}_p - m_B} \gamma_5 \not{p}_{\pi^-} - \gamma_5 \not{p}_{K^-} \frac{\mathcal{H}_{\mathcal{P}N}/3}{\not{p}_{\pi^-} + \not{p}_p - m_{\tilde{N}}} \Gamma_P \not{p}_{\pi^-} \right) u_p. \qquad (11)$$

We show the total cross section, rescaled by the factor Γ_Θ/MeV, as functions of pion energy E_π at the proton rest frame in Fig. 2 depending on the Θ^+ parity, and varied $-4 \leqslant \mathcal{H}_{\mathcal{P}N} \leqslant 4$. Note that the sign of $\mathcal{H}_{\mathcal{P}N}$ is very important. If $\mathcal{H}_{\mathcal{P}N} > 0$ (< 0), two contributions will have constructive (destructive) interference. Thus our results differ from the previous results where only the n contribution was included. Also the cross section is sensitive to $\mathcal{H}_{\mathcal{P}N}$, and may be useful to fix $\mathcal{H}_{\mathcal{P}N}$. Following the spirit of chiral perturbation theory at lowest order, we did not include model-dependent form factors, keeping in mind that our results get unreliable as the E_π gets as large as ~ 2.5 GeV.

Note that the even parity and the odd parity cases can be distinguished from the cross section for the parity-odd case is smaller than that for the parity-even case.

4. Photo-production of Θ^+

In order to study photo-production of Θ^+ on nucleons, we need to know the magnetic dipole interaction terms. For the nucleon octet,

$$\mathcal{L} = -\frac{e}{4m_B} \text{Tr}\left[\bar{B} \sigma_{\mu\nu} F^{\mu\nu} \left(\kappa_D \{Q, B\} + \kappa_F [Q, B]\right)\right].$$

The anomalous magnetic moments of nucleons are

$$\kappa_p = \kappa_F + \frac{1}{3}\kappa_D, \qquad \kappa_n = -\frac{2}{3}\kappa_D,$$

at tree-level chiral Lagrangian. Using $\kappa_p = 1.79$ and $\kappa_n = -1.91$, we get $\kappa_D = 2.87$ and $\kappa_F = 0.836$. For the pentaquark baryon $\mathcal{P}$, the relevant term is

$$-\frac{e\kappa_{\mathcal{P}}}{4m_{\mathcal{P}}} q_i \bar{\mathcal{P}}_i \sigma_{\mu\nu} F^{\mu\nu} \mathcal{P}_i \to -\frac{e\kappa_\Theta}{4m_\Theta} \bar{\Theta}^+ \sigma_{\mu\nu} F^{\mu\nu} \Theta^+.$$

We expect that $|\kappa_\Theta (\equiv \kappa_{\mathcal{P}})| \approx |\kappa_D| \approx |\kappa_F|$. On the other hand, a calculation in soliton picture predicts that $\kappa_\Theta \approx 0.3$, which is rather small [13]. We vary κ_Θ between -1 and 1. We ignore transition magnetic moments between nucleon octet and pentaquark antidecuplet, since this transition involves qqq and $qqqq\bar{q}$.

The relevant Feynman diagrams for $\gamma n \to K^- \Theta^+$ are shown in Fig. 3. One salient feature of our approach based on chiral perturbation theory is the

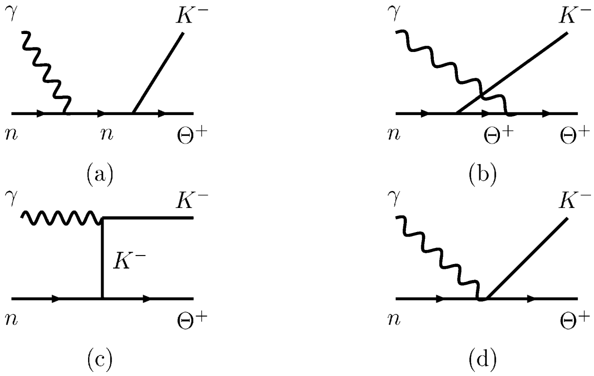

Fig. 3. Feynman diagrams for $\gamma n \to K^- \Theta^+$.

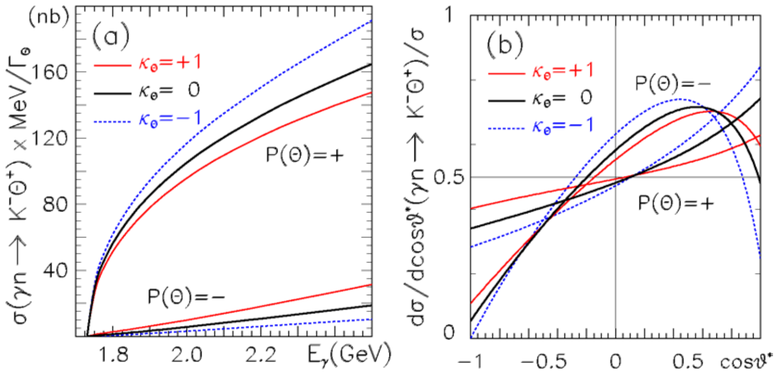

Fig. 4. (a) Cross sections for $\gamma n \to K^- \Theta^+$ and (b) the angular distribution for $E_\gamma = 2$ GeV in the center of momentum frame.

existence of a contact term for $\gamma K^- n \Theta^+$ vertex (Fig. 3(d)) that arises from the $\mathcal{C}_{\mathcal{P}N}$ term in Eq. (4c) with Eq. (2c), which is necessary to recover $\mathrm{U}(1)_{\mathrm{em}}$ gauge invariance within spontaneously broken global chiral symmetries. The resulting amplitude is

$$\mathcal{M} = \frac{e\mathcal{C}_{\mathcal{P}N}}{\sqrt{2}f} \epsilon_\mu \bar{u}_{\Theta^+} F^\mu u_n, \tag{12a}$$

$$\begin{aligned} F^\mu = {} & -\Gamma_P \not{p}_K \frac{\kappa_n}{4m_B} \frac{1}{\not{p}_\gamma + \not{p}_n - m_B} [\gamma^\mu, \not{p}_\gamma] + \Gamma_P \gamma^\mu \\ & + \left(\gamma^\mu - \frac{\kappa_\Theta}{4m_\Theta} [\gamma^\mu, \not{p}_\gamma] \right) \frac{1}{\not{p}_n - \not{p}_K - m_\Theta} \\ & \times \Gamma_P \not{p}_K \\ & - \frac{(2p_K - p_\gamma)^\mu}{(p_K - p_\gamma)^2 - m_K^2} \Gamma_P (\not{p}_K - \not{p}_\gamma). \end{aligned} \tag{12b}$$

The cross sections and the angular distributions in the center of momentum frame are shown in Fig. 4(a) and (b). Note that the parity-even case has larger cross section, and has a sharp rise near the threshold. The angular distribution shows that the forward/backward scattering is suppressed in the negative parity case, whereas the forward peak is present in the positive parity case. Therefore the angular distribution could be another useful tool to determine the parity of Θ^+. Therefore, once $\mathcal{C}_{\mathcal{P}N}^2$ is determined from Γ_Θ, one could determine the parity of Θ^+, and make a rough estimate of κ_Θ from the photo-production cross section.

5. Including light vector mesons

One can also introduce light vector mesons ρ_μ, which transforms as

$$\rho_\mu(x) \to U(x)\rho_\mu(x)U^\dagger(x) + U(x)\partial_\mu U^\dagger(x), \tag{13}$$

under global chiral transformations [14]. Then $\rho_\mu(x)$ transforms as a gauge field under local SU(3)'s defined by Eq. (1), as V_μ does. The covariant derivative $\mathcal{D}_\mu$ can be defined using ρ_μ instead of V_μ. Note that $(\rho_\mu - V_\mu)$ has a simple transformation property under chiral transformation

$$(\rho_\mu - V_\mu) \to U(x)(\rho_\mu - V_\mu)U^\dagger(x),$$

and it is straightforward to construct chiral invariant Lagrangian using this new field. In terms of a field strength tensor $\rho_{\mu\nu}$,

$$\mathcal{L}_\rho = -\frac{1}{2}\operatorname{Tr}\bigl(\rho_{\mu\nu}\rho^{\mu\nu}\bigr) + \frac{1}{2}m_\rho^2 \operatorname{Tr}(\rho_\mu - V_\mu)^2 + \alpha\bigl[\bar{\mathcal{P}}(\not\rho - \not V)B + \bar{B}(\not\rho - \not V)\mathcal{P}\bigr] + \cdots.$$

It is important to notice that $N\Theta^+K^*$ coupling should be highly suppressed, since it can appear only in combination of $(\rho_\mu - V_\mu)$, which vanishes in the low-energy limit. In other words, the low-energy theorem is violated if one includes only $n\Theta^+K^*$ diagram, without including the $n\Theta^+K\pi$ contact term arising from the $\not V$ term. Therefore, one should be cautious about claiming that the K^* exchange is important in $\pi^- p \to K^-\Theta^+$. Detailed numerical analysis of vector-meson exchange is straightforward, but beyond the scope of the present work and will be pursued elsewhere [15].

6. Conclusion

In conclusion, we constructed a chiral Lagrangian involving pentaquark baryon antidecuplet and octet, the ordinary nucleon octet and Goldstone bosons. Using this Lagrangian, we derived the Gell-Mann–Okubo formula and the mixing between the pentaquark antidecuplet and pentaquark octet. We also calculated the cross sections for $\pi^- p \to K^-\Theta^+$ and $\gamma n \to K^-\Theta^+$ for $J^P = \frac{1}{2}^\pm$. In particular, we emphasized that it is very important to respect chiral symmetry properly in order to get correct amplitudes for these processes. All these observables depend on parameters in our chiral Lagrangian, which have relations with underlying QCD, but are incalculable from QCD at present. Photo-production data will be particularly useful in identifying the parity of Θ^+, because the threshold behavior of the cross section and its angular distributions strongly depend on the parity. Once the coupling $\mathcal{C}_{\mathcal{P}N}$ is determined from the decay width of Θ^+, then the parity and other couplings $\mathcal{H}_{\mathcal{P}N}$ and κ_Θ could be determined from the hadro/photo-production cross sections for Θ^+. It is straightforward to apply our approach to other related processes such as $\gamma p \to K^0\Theta^+$ or $K^+ p \to \pi^+\Theta^+$, or other pentaquark baryons. Finally we have outlined how to incorporate the vector meson degrees of freedom in our scheme, the details of which will be discussed in the separate publication [15].

Acknowledgements

This work is supported in part by KOSEF through CHEP at Kyungpook National University and by the BK21 program. The research of J.L. in the High Energy Physics Division at Argonne National Laboratory is supported by the US Department of Energy, Division of High Energy Physics, under Contract W-31-109-ENG-38. J.L. thanks KAIST and KIAS for their hospitality during this work. T.L. is supported in part by BK21 Program and Korea Research Foundation Grant (KRF-2004-015-C00095).

References

[1] LEPS Collaboration, T. Nakano, et al., Phys. Rev. Lett. 91 (2003) 012002.
[2] DIANA Collaboration, V.V. Barmin, et al., Phys. At. Nucl. 66 (2003) 1715, Yad. Fiz. 66 (2003) 1763 (in Russian).
[3] CLAS Collaboration, S. Stepanyan, et al., hep-ex/0307018.
[4] CLAS Collaboration, V. Kubarovsky, S. Stepanyan, hep-ex/0307088.
[5] SAPHIR Collaboration, J. Barth, et al., Phys. Lett. B 572 (2003) 127.
[6] A.E. Asratyan, A.G. Dolgolenko, M.A. Kubantsev, hep-ex/0309042.
[7] D. Diakonov, V. Petrov, M.V. Polyakov, Z. Phys. A 359 (1997) 305, hep-ph/9703373.
[8] A.V. Manohar, hep-ph/9606222.
[9] R.L. Jaffe, F. Wilczek, hep-ph/0307341.
[10] D. Diakonov, V. Petrov, hep-ph/0310212.
[11] Y.S. Oh, H.C. Kim, S.H. Lee, hep-ph/0310117.
[12] R.N. Cahn, G.H. Trilling, hep-ph/0311245.
[13] H.C. Kim, hep-ph/0308242.
[14] M. Bando, T. Kugo, K. Yamawaki, Nucl. Phys. B 259 (1985) 493.
[15] Work in progress.