Two topics in chiral perturbation theory: $X(3872)$ and pentaquark baryons

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I discuss two applications of chiral lagrangian to multiquark hadrons : (i) the dipion invariant mass spectrum in $X(3872) \to J/\psi\pi\pi$ and (ii) some properties of pentaquark baryons.

I. DIPION SPECTRUM IN $X(3872) \to J/\psi\pi\pi$

B decay is a good place to look for charmonium states which are either above the $D^{(*)}\bar{D}^{(*)}$ threshold or have quantum numbers that are not accessible from n^3S_1 states by cascade decays. $2^3P_J, ^3D_J$ and 1^1P_1 states are such examples [1]. Recently, Belle collaboration reported a new narrow resonance $X(3872)$ in $B \to XK \to (J/\psi\pi\pi)K$ decay channel [2].Following this observation, there appearted a lot of attempt to understand this new narrow resonance (see Ref. [3] for a recent review). In particular, Pakvasa and Suzuki pinned down possible quantum numbers for $J^{PC}(X)$ [4] as follows:

- If it is a charmonium, it should be either $1^3D_2(2^{--})$ or $2^1P_1(1^{+-})$.
- If $X(3872)$ is the $D\overline{D^*}$ molecular state, it should be either $J^{PC} = 1^{+-}$ with $I = 0$, or $J^{PC} = 1^{++}$ with $I = 1$.

The angular distribution of dipions would be useful to determine the J^{PC} quantum number of $X(3872)$ [4].

In this talk, I argue that the dipion invariant mass spectrum in $X \to J/\psi\pi\pi$ provides independent information on the nature of $X(3872)$ (see Ref. [5] for details). Our method is complementary to the angular correlations suggested in Ref. [4], and already seem to eliminate a possibility that $X = {}^1P_1$ state.

Hadronic transition between heavy quarkonia can be described in terms of QCD multipole expansion or chiral perturbation theory. Since the allowed dipion invariant mass in $X(k, \epsilon_X) \to J/\psi(k^{'}, \epsilon_\psi)\ \pi(p_1)\pi(p_2)$ is

$$2m_\pi \le m_{\pi\pi} \le (M_X - M_{J/\psi}) = 775 \text{ GeV } \approx m_\rho,$$

both approaches will be (marginally) suitable to our purpose.

Under global chiral $SU(3)_L \times SU(3)_R$, the pion field $\Sigma(x) \equiv exp(2i\pi(x)/f_\pi)$ transforms nonlinearly as

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

Under parity P and charge conjugation C, the pion fields transform as

$$\begin{aligned} P: &\ \pi(t,\vec{x}) \to -\pi(t,-\vec{x}),\ \ \Sigma(t,\vec{x}) \to \Sigma^\dagger(t,-\vec{x}) \\ C: &\ \ \pi(t,\vec{x}) \to \pi(t,\vec{x})^T, \quad \Sigma(t,\vec{x}) \to \Sigma(t,\vec{x})^T \end{aligned} \qquad (1)$$

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TABLE I: Transformation properties of X_v, ψ_v and v under parity and charge conjugation.

Fields	P	C
v^μ	$v_\mu = (v^0, -\vec{v})$	$v^\mu = (v^0, \vec{v})$
ψ_v^μ	$\psi_{v\mu}$	$-\psi_v^\mu$
$X_v^{\mu\nu}(^3D_2)$	$-X_{v\mu\nu}$	$-X_v^{\mu\nu}$
$X_v^\mu(^1P_1)$	$-X_{v\mu}$	$-X_v^\mu$
$X_v^\mu(J^{PC} = 1^{++})$	$-X_{v\mu}$	X_v^μ

We impose parity P and charge conjugation C as symmetries of the effective chiral lagrangian, since all the decays under consideration occur through strong interactions.

Since the final charmonium is moving very slowly in the rest frame of the initial state $X(3872)$, we can use the heavy particle effective theory approach, by introducing a velocity dependent field $X_v(x) \equiv Xe^{im_X v\cdot x}$, and similarly for J/ψ field $\psi_v(x)$ [6]. Then we can construct chiral lagrangian using the pion field $\Sigma(x)$ and chiral singlets $X_v(x)$, $\psi_v(x)$, $\epsilon_{\mu\nu\alpha\beta}$, and the velovity v_μ. Transformation properties of $X_v(x)$, $\psi_v(x)$ and v^μ under parity and charge conjugation are given in Table 1.

A remark is in order before we procees. For $\psi^{'} \to J/\psi\pi\pi$ or $\Upsilon(3S) \to \Upsilon(1S)\pi\pi$, there are 4 independent chirally invariant operators to lowest order in pion momenta:

$$\begin{aligned} \mathcal{M}\ &\sim\ \epsilon\cdot\epsilon^{'}\left[q^2 + BE_1E_2 + Cm_\pi^2\right] \\ &+\ D\left(\epsilon\cdot p_1\epsilon^{'}\cdot p_2 + \epsilon\cdot p_2\epsilon^{'}\cdot p_1\right), \end{aligned}$$

with $B, C, D \sim O(1)$. QCD multipole expansion predicts that $B = D = 0$, which is consistent with $\psi^{'} \to J/\psi\pi\pi$ or $\Upsilon(2S) \to \Upsilon(1S)\pi\pi$ decays, but not with $\Upsilon(3S) \to \Upsilon(1S)\pi\pi$. In chiral lagrangian approach, B, C, D being apriori unknown with similar sizes of $O(1)$. One has to fix B, C, D from the $m_{\pi\pi}$ spectrum and predict other observables [7]. Namely, one cannot make definite predictions for the $m_{\pi\pi}$ spectrum in $^3S_1 \to ^3S_1\pi\pi$. On the other hand, there is only single chirally invaraint operator to lowest order in pion momenta in the three options for the $J^{PC}(X)$ we consider in this talk, and we can make a definite prediction for $m_{\pi\pi}$ sprectrum.

If $X(3872) = 1^3D_2(2^{--})$, then the decay $X \to J/\psi\pi\pi$ is described by the following chiral lagrangian:

$$\mathcal{L} = g(^3D_2)\ \epsilon^{\mu\nu\alpha\beta}\ v_\mu\psi_{v,\nu}X_{v,\alpha\rho}\ \text{Tr}\ \left[\partial_\beta U\partial^\rho U^\dagger\right] + h.c. \quad (2)$$

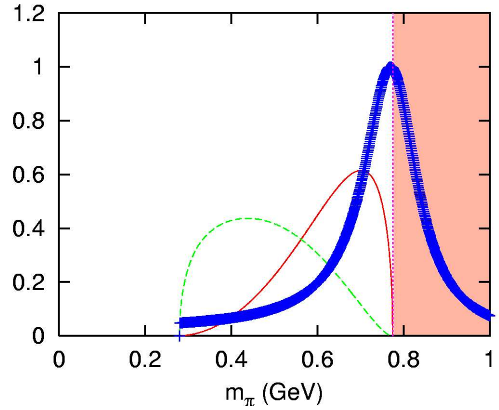

FIG. 1: Dipion invariant mass spectra for $J^{PC}(X) = 2^{--}$ (red), $J^{PC}(X) = 1^{+-}$ (green) and $J^{PC}(X) = 1^{++}$ (blue). The vertical scale is arbitrary for each case, and the area under a curve has no meaning in this plot. The shaded (pink) region is kinematically forbidden region.

where the coupling $g(^3D_2)$ is unknown parameter, that should be determined by the data or could be calculated within QCD multipole expansion. Note that there is only one operator that contributes to the decay $1^3D_2(2^{--}) \to J/\psi\pi\pi$, so that we can predict the $\pi\pi$ spectrum without any ambiguity. Now it is straightforward to calculate the dipion invariant mass spectrum using the above lagrangian. The resulting spectrum is shown in Fig. 1 in the (red) solid curve. Note that the dipion invariant mass has a peak at high $m_{\pi\pi}$ region, which is consistent with the preliminary Belle data [2].

If X is a charmonium $^1P_1$ state with $J^{PC} = 1^{+-}$ and $I = 0$, the relevant chiral lagrangian is given by

$$\mathcal{L} = g(^1P_1)\ \epsilon^{\mu\nu\alpha\beta}\ X_{v\mu}\psi_{v\nu}\ \mathrm{Tr}\ \partial_\alpha U \partial_\beta U^\dagger \qquad (3)$$

where the coupling $g(^1P_1)$ is an unknown parameter. If X is a $D\overline{D^*}$ molecular state with $J^{PC} = 1^{+-}$, we cannot apply QCD multipole expansion. Still the chiral lagrangian shown above is applicable.

We show the resulting $\pi\pi$ invariant mass spectrum in Fig. 1 in the (green) dotted curve. Note that the dipion spectrum has a peak at low $m_{\pi\pi}$ region, if X has $J^{PC} = 1^{+-}$, which is disfavored by the preliminary Belle data [2]. Once high statistics data is obtained, one can easily check if $J^{PC}(X) = 1^{+-}$ is a correct assignment or not.

Let us finally discuss the case $X(3872) = (1^{++}, I = 1)$. This case includes that the decaying state is $I = 1$ $D\overline{D^*}$ molecular state and the final dipion is in $I = 1$, which would be dominated by ρ meson. Since $\rho^0 \to \pi^0\pi^0$ is forbidden by angular momentum conservation and Bose symmetry, the $\pi\pi$ in $X \to J/\psi\pi\pi$ should be charged pions in this case. The $m_{\pi\pi}$ spectrum is basically the Breit–Wigner profile of the ρ resonance, which is consistent with the current Belle data. In this case, the polarization of ρ and J/ψ tends to be perpendicular with each other, which can be tested by measuring the three–momentum of a lepton in $J/\psi \to l^+l^-$ and the three–momentum of a pion in $\rho \to \pi\pi$ decay.

Summarizing the first part of may talk, I pointed out that the dipion invariant mass spectrum in $X(3872) \to J/\psi\pi\pi$ could be useful in determination of J^{PC} quantum number of the newly observed resonance $X(3872)$. In particular, the current preliminary data seems to already exclude the possibility $X = {}^1P_1$, since the dipion invariant mass spectrum has a peak at low $m_{\pi\pi}$ region in this case. This is apparently disfavored by the preliminary Belle data.

II. PENTAQUARK BARYONS

Recently, five independent experiments reported observations of a new baryonic state $\Theta^+(1540)$ with a very narrow width < 5 MeV [8], which is likely to be a pentaquark state ($uudd\bar{s}$) predicted by Diakonov et al. [9]. 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. In the following, 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 antidecuplets, 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.

In the previous section, we already introduced Goldstone boson fields including pions and their transformation property under chiral $SU(3)$. Let us denote baryon octet including nucleons by B, and antidecuplet including Θ^+ by $\mathcal{P}$. It is convenient to define another field $\xi(x)$ by $\Sigma(x) \equiv \xi^2(x)$, which transforms as [10]

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

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:

$$\begin{aligned} V_\mu &= \tfrac{1}{2}(\xi^\dagger\partial_\mu\xi + \xi\partial_\mu\xi^\dagger), \quad V_\mu \to UV_\mu U^\dagger + U\partial_\mu U^\dagger, \\ A_\mu &= \tfrac{i}{2}(\xi^\dagger\partial_\mu\xi - \xi\partial_\mu\xi^\dagger), \quad A_\mu \to UA_\mu U^\dagger. \end{aligned} \qquad (4)$$

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

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

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

$$\mathcal{D}_\mu B = \partial_\mu B + [V_\mu, B].$$

Chiral symmetry is explicitly broken by non-vanishing 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 $F_{\mu\nu} = \partial_\mu \mathcal{A}_\nu - \partial_\nu \mathcal{A}_\mu$:

$$\partial_\mu \Sigma \;\to\; \mathcal{D}_\mu \Sigma \equiv \partial_\mu \Sigma + ie\mathcal{A}_\mu [Q, \Sigma], \qquad (5a)$$

$$V_\mu \;\to\; V_\mu + \tfrac{ie}{2}\mathcal{A}_\mu(\xi^\dagger Q \xi + \xi Q \xi^\dagger), \qquad (5b)$$

$$A_\mu \;\to\; A_\mu - \tfrac{e}{2}\mathcal{A}_\mu(\xi^\dagger Q \xi - \xi Q \xi^\dagger), \qquad (5c)$$

where $Q \equiv$ 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

$$\mathcal{L} = \mathcal{L}_\Sigma + \mathcal{L}_B + \mathcal{L}_\mathcal{P}, \qquad (6)$$

where

$$\mathcal{L}_\Sigma \;=\; \frac{f_\pi^2}{4}\,\mathrm{Tr}\big[\mathcal{D}_\mu \Sigma^\dagger \mathcal{D}^\mu \Sigma - 2\mu m(\Sigma + \Sigma^\dagger)\big], \qquad (7a)$$

$$\mathcal{L}_B \;=\; \mathrm{Tr}\,\overline{B}(i\not{\mathcal{D}} - m_B)B + D\,\mathrm{Tr}\,\overline{B}\gamma_5\{\not{A}, B\} + F\,\mathrm{Tr}\,\overline{B}\gamma_5[\not{A}, B], \qquad (7b)$$

$$\mathcal{L}_\mathcal{P} \;=\; \overline{\mathcal{P}}(i\not{\mathcal{D}} - m_\mathcal{P})\mathcal{P} + \mathcal{C}_{\mathcal{P}N}\big(\overline{\mathcal{P}}\Gamma_P \not{A} B + \overline{B}\Gamma_P \not{A}\mathcal{P}\big) + \mathcal{H}_{\mathcal{P}N}\overline{\mathcal{P}}\gamma_5\not{A}\mathcal{P}, \qquad (7c)$$

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

$$\mathcal{L}_m = \alpha_m\,\overline{\mathcal{P}}(\xi m \xi + \xi^\dagger m \xi^\dagger)\mathcal{P}. \qquad (8)$$

Expanding this, we get the mass splittings $\Delta m_i \equiv m_i - m_\mathcal{P}$ within the antidecuplet. If the newly observed state at a mass 1862 ± 2 MeV is identified as $\Xi_{3/2}$, we find

$$m_{\widetilde{N}} = 1647\,\mathrm{MeV}, \quad m_{\widetilde{\Sigma}} = 1755\,\mathrm{MeV}. \qquad (9)$$

One can consider a mixing between pentaquark antidecuplet $\mathcal{P}_{abc}$ and pentaquark octet $\mathcal{O}^a{}_b$ or with the ordinary baryon octet in a similar fashion [12]. But we do not pursue it furthermore here, since the current data on baryon sectors are not enough to study such mixings in details.

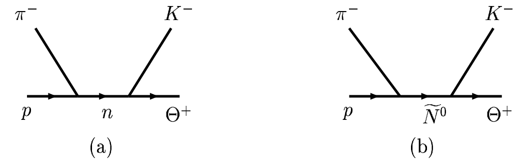

FIG. 2: Feynman diagrams for $\pi^- p \to K^- \Theta^+$.

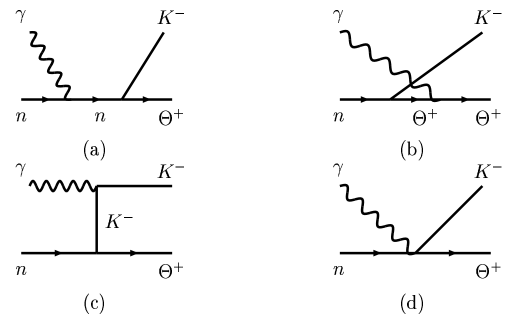

FIG. 3: Feynman diagrams for $\gamma n \to K^- \Theta^+$.

Parameters in the above lagrangian are taken to have the following numerical values: $m_B \approx 940\,\mathrm{MeV}$ 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 $K^0 p$ modes:

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

The narrow width of Θ implies $\mathcal{C}^2_{\mathcal{P}N} \sim 10^{-3}$. Understanding such a small $\mathcal{C}_{\mathcal{P}N}$ in strong interaction is one of the important issues in pentaquark physics.

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 $\mathcal{H}_{\mathcal{P}N} \sim O(1)$, without any suppression as in $\mathcal{C}_{\mathcal{P}N}$. This fact makes it impossible to predict the cross section for $\pi^- p \to K^- \Theta^+$, even if we ignore the uncertainties from form factors. In Figs. 2 (a) and (b), we show the relevant Feynman diagrams. Note that only Fig. 2 (a) was considered in the literature. However, there is an $s-$channel $\widetilde{N}^0(1647)$ exchange diagram [Fig. 2 (b)] in our chiral lagrangian, since Θ^+ is not an SU(3) singlet, but belongs to the antidecuplet. Therefore one has to keep both Figs. 2 (a) and (b) in order to get an amplitude with correct SU(3) flavor symmetry. Since the coupling $\mathcal{H}_{\mathcal{P}N}$ is not known, and we cannot make a clean prediction for $\pi^- p \to K^- \Theta^+$ cross section.

Next, let us discuss photoproduction of Θ^+ on nucleons. For this purpose, we included the anomalous magnetic moments for nucleon octet (κ_n and κ_p) and pentaquark baryon Θ (κ_Θ). We expect that $|\kappa_\Theta| \sim |\kappa_n| \sim |\kappa_p|$. The relevant Feynman diagrams for $\gamma n \to K^- \Theta^+$

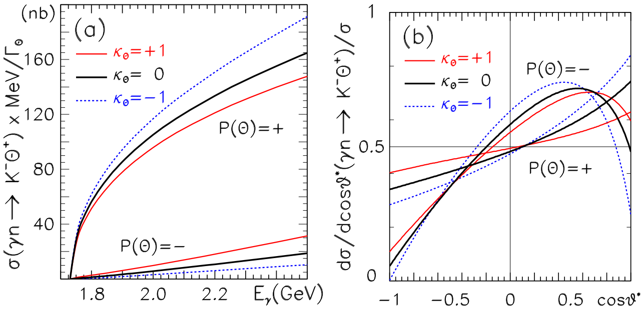

FIG. 4: (a) Cross sections for $\gamma n \to K^- \Theta^+$ and (b) the angular distribution for $E_\gamma = 2$ GeV in the center of momentum frame.

are shown in Fig. 3. One salient feature of our approach based on chiral perturbation theory is the 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. (7c) with Eq. (5c), which is necessary to recover $\mathrm{U}(1)_{\mathrm{em}}$ gauge invariance within spontaneously broken global chiral symmetries.

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 $\mathcal{C}^2_{\mathcal{P}N}$ is determined from Γ_Θ, one could determine the parity of Θ^+, and make a rough estimate of κ_Θ from the photoproduction cross section. Again we have to keep in mind that there is additional model dependence from unknown form factors.

One can also introduce light vector mesons ρ_μ, which transforms as $\rho_\mu(x) \to U(x)\rho_\mu(x)U^\dagger(x) + U(x)\partial_\mu U^\dagger(x)$, under global chiral transformations [13]. 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 $\rho_{\mu\nu}$,

$$\begin{aligned}\mathcal{L}_\rho &= -\frac{1}{2}\,\mathrm{Tr}(\rho_{\mu\nu}\rho^{\mu\nu}) + \frac{1}{2}\,m_\rho^2\,\mathrm{Tr}(\rho_\mu - V_\mu)^2 \\ &+ \alpha\left[\overline{\mathcal{P}}(\not\rho - \not V)B + \overline{B}(\not\rho - \not V)\mathcal{P}\right] + \ldots\end{aligned}$$

It is important to notice that $N\Theta^+K^*$ coupling ($\propto \alpha$) 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 $\pi^- p \to K^- \Theta^+$.

Summarizing the second half of my talk, 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 discussed $\pi^- p \to K^- \Theta^+$ 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.

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