$K^*$-couplings for the antidecuplet excitation

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Abstract

We estimate the coupling of the $K^*$ vector meson to the $N \to \Theta^+$ transition employing unitary symmetry, vector meson dominance, and results from the GRAAL Collaboration for $\eta$ photoproduction off the neutron. Our small numerical value for the coupling constant is consistent with the non-observation of the $\Theta^+$ in recent CLAS searches for its photoproduction. We also estimate the $K^*$-coupling for the $N \to \Sigma^*$ excitation, with $\Sigma^*$ being the $\Sigma$-like antidecuplet partner of the $\Theta^+$-baryon.

1 Introduction

The experimental status of the exotic baryon $\Theta^+$ is rather uncertain. Various collaborations have given either positive results for its observation, or null results, casting doubts on the existence of the $\Theta^+$ and on the correctness of the positive claims (see recent experimental reviews [1, 2]). The apparently most impressive searches today are the recent CLAS Collaboration publications [3, 4, 5, 6] with null results for $\Theta^+$-photoproduction in several final-state channels. Nevertheless, investigations are continuing, and several new (“after-CLAS”) dedicated experiments have given both positive [7, 8, 9, 10] and null [11, 12, 13, 14, 15] results.

The situation is complicated by the fact that there are no data sets, from independent groups, with exactly overlapping conditions (initial and/or final states, kinematical regions). Therefore, when using any null result of one group to reject the positive result of another, we need to apply some theoretical models/assumptions. Here we also encounter numerous uncertainties. For instance, current approaches to the theoretical description of $\Theta$-production (for example, its photoproduction off the nucleon [16, 17, 18, 19, 20, 21, 22]) are mainly based on $K$ and $K^*$ meson or reggeon exchanges. The vertex for the $KN\Theta$-coupling may be considered as known, if one assumes the spin-parity and width of $\Theta^+$ to be known (the
corresponding form factor is still a problem, of course). Contrary to this, properties of the $K^*$ exchange are totally unknown. However, they may be essential, e.g., for comparisons of $\Theta^+$-photoproduction off the proton and/or the neutron.

In the present note, we obtain at least rough estimates of the $K^*N\Theta$-coupling. Similar estimates will be provided for $K^*$-vertices contributing to the excitation $N \rightarrow \Sigma^*$. To derive the present results, we relate these vertices to the radiative vertex, extracted recently \cite{23} from preliminary experimental evidence on $N(1675)$, obtained by the GRAAL Collaboration \cite{24}, with preliminary confirmation from the CB-ELSA-TAPS \cite{25} and LNS \cite{26} Collaborations.

2 Transition vertices

Let us consider transitions between two baryons $B_1 \rightarrow B_2$ through radiation of a photon or $K^*$. As a first step, we may assume the flavor symmetry $SU(3)_F$ to be exact. Then, the corresponding vertices for the vector meson octet are generated in QCD by the quark vector current $J^a_\mu$. This is an octet in $SU(3)_F$, and the superscript $a$ specifies its components. The $K^*$ meson, or any other member of the vector meson octet, is coupled to a particular combination of the components of the $J^a_\mu$. The electromagnetic current is also related (proportional) to one of those components. Since it is conserved, and the $SU(3)_F$ is assumed to be exact, all other components should be conserved as well. If both initial and final baryons have $J^P = 1/2^+$, then the transition vertex $V^a_\mu(B_2B_1)$ has the general form

$$V^a_\mu(B_2B_1) = \langle B_2|J^a_\mu|B_1\rangle = \langle B_2|\left( f_1^a(B_1B_2)\gamma_\mu + f_2^a(B_1B_2)i\sigma_{\mu\nu}q^\nu m_1 + m_2 + f_3^a(B_1B_2)\frac{q_\mu}{m_1 + m_2} \right)|B_1\rangle, \quad (1)$$

where $q = p_2 - p_1$ is the momentum transfer, $m_1$ and $m_2$ are the masses of the corresponding baryons, and the form factors $f_i^a(B_1B_2)$ are invariant functions of $q^2$. For a specific vector meson $M$, we obtain the corresponding vertex $MB_1B_2$ of the same structure, with coefficients $f_i(MB_1B_2)$ which are specific combinations of $f_i^a(B_1B_2)$.

Conservation of the current $J^a_\mu$ implies that

$$f_1^a(B_1B_2) \cdot (m_2 - m_1) = -f_3^a(B_1B_2) \frac{q^2}{m_1 + m_2}. \quad (2)$$

In the present note, we are interested in transitions between baryons of different flavor multiplets and, hence, of different masses. For such transitions, $f_1^a$ is proportional to $q^2$ and, thus, disappears at $q^2 = 0$, in contrast to $f_3^a$. For the case of exact $SU(3)_F$, the last term of the vertex (1) does not provide any non-vanishing physically meaningful contribution (in total analogy with a similar term in the photon vertex).

Though the situation might change after accounting for the violation of $SU(3)_F$, we advocate the tensor coupling (i.e., the second term of Eq. (1)) to have the leading role at small $q^2$ for non-diagonal transitions. The relative contributions of other couplings should be suppressed. For the radiative decay $B_1 \rightarrow B_2 \gamma$ (or $B_2 \rightarrow B_1\gamma$), with the real photon having $q^2 = 0$, only the second term in the vertex is physically meaningful, and

$$f_2(\gamma B_1 B_2)|_{q^2=0} = \mu(B_1 \rightarrow B_2) \cdot (m_1 + m_2), \quad (2)$$

where $\mu(B_1 \rightarrow B_2)$ is the transition magnetic moment.
The dominance of the tensor coupling over the vector one in hadronic non-diagonal transitions (between different flavor multiplets) is supported by the analysis \cite{27} of data for transitions between 10 and 8 multiplets. For diagonal transitions $B \rightarrow B$, the conservation of the current $J_{\mu}^B$ does not require $f_1^B(BB)$ to vanish at $q^2 \rightarrow 0$, though the “tensor dominance” may still be true as well, at least for some components of the current. For example, experimental data compilation \cite{28} suggests that for the vertex $\rho NN$ the ratio $f_2/f_1$ is large ($\sim 6$, the tensor-over-vector dominance is true), but for the vertices $\omega NN$ and $\phi NN$ it is small ($< 0.2$, the dominance is untrue).

Regrettably, many phenomenological calculations of $\Theta$-production (e.g., \cite{16,17,18,19,20}) have used only the vector coupling for vector-meson exchanges, assuming a non-vanishing $f_1^\Lambda|_{q^2 \rightarrow 0}$ and neglecting $f_2^\Lambda(q^2)$. To our best knowledge, dominance of the tensor coupling has been used only in Refs. \cite{21,22}.

3 $SU(3)_F$ relations

Let us consider the tensor couplings for the transitions $8 \rightarrow \overline{10}$. An octet baryon cannot be transformed into an antidecuplet member using a unitary singlet meson, while it can be done through coupling with a unitary octet meson. Since the product $8 \otimes 8$ contains $\overline{10}$ only once, all the appearing coupling constants may be expressed through one of them. The corresponding relations may be found by means of the Clebsch-Gordan coefficients for the group $SU(3)$ \cite{29}. This method was used in Ref. \cite{30} to find relations between coupling constants for transitions $\overline{10} \rightarrow 8$ with radiation of the octet pseudoscalar mesons. For our case (vector mesons), these results would need some modification, because the vector mesons have an essential singlet-octet mixing, absent for the pseudoscalar mesons. Instead, we will apply here another (equivalent) way which uses only the standard (and more widely familiar) $SU(2)$-formalism of the Clebsch-Gordan coefficients, with the corresponding phase conventions.

Our approach is based on the fact that the group $SU(3)_F$ has three $SU(2)$-subgroups \cite{29}: the familiar isotopic spin ($I$-spin) group $SU(2)_I$, with the spinor $(u, d)$ and singlet $s$; the $U$-spin group $SU(2)_U$, with the spinor $(d, s)$ and singlet $u$; the $V$-spin group $SU(2)_V$, with the spinor $(s, u)$ and singlet $d$. In the framework of exact $SU(3)_F$, the photon corresponds to the $U$-singlet component of octet.

It is sufficient for our present purpose to use any two of the three subgroups. Most convenient for us are the isotopic spin and $U$-spin subgroups (note that all members of any $U$-spin multiplet have the same electric charge).

With sufficient accuracy for our purpose here, we assume the nonet of ground state vector mesons to have ideal octet-singlet mixing, such that $\phi = s\bar{s}$, while the $\omega$-meson contains light quarks only, $\omega = (u\bar{u} + d\bar{d})/\sqrt{2}$. The ground state baryons are taken, as usual, to be members of an unmixed octet. For the antidecuplet, we neglect possible mixings, considering them to be parametrically small \cite{31}.

In terms of $U$-spin, the proton belongs to the doublet $(p, \Sigma^+)$, with $U = 1/2$, while the proton-like member of the antidecuplet, $p^*$, enters the $U$-spin quartet $(\Theta^+, p^*, \Sigma^{*+}, \Xi^{*+}_{3/2})$, with $U = 3/2$. The vector meson nonet contains two independent $U$-singlet combinations, which cannot change $U$-spin and, thus, cannot transform $p \rightarrow p^*$. We may take them as $\rho^0 + \omega = \sqrt{2}u\bar{u}$ and $\omega + \phi/\sqrt{2} = (u\bar{u} + d\bar{d} + s\bar{s})/\sqrt{2}$ (note that the $U$-singlet component of an octet may change the isotopic spin, as does the photon).

With the above $U$-singlet combinations, we immediately obtain two relations

$$f_2(\rho^0 pp^*) = -f_2(\omega pp^*) = \frac{1}{\sqrt{2}} f_2(\phi pp^*) .$$ (3)
The meson $K^* = d\bar{s}$ has $U = 1$ with $U_3 = +1$. It is a member of a $U$-spin triplet, other members of which are the combination $(\omega - \rho^0)/2 - \phi/\sqrt{2} = (d\bar{d} - s\bar{s})/\sqrt{2}$, with $U_3 = 0$, and $(\bar{K}^*0)$, with $U_3 = -1$. When applying to this triplet the standard $SU(2)$ Clebsch-Gordan coefficients, which couple doublet ($J = 1/2$) and triplet ($J = 1$) into quartet ($J = 3/2$), and accounting for relations (3), we obtain a simple new relation

$$f_2(K^*0 p \Theta^+) = -\sqrt{6} f_2(\rho^0 p p^*).$$

(4)

It is then easy to use the standard isotopic ($I$-spin) relations and obtain the neutron couplings

$$f_2(K^{*+} n \Theta^+) = -f_2(K^*0 p \Theta^+), \quad f_2(\rho^0 n n^*) = -f_2(\rho^0 p p^*), \quad f_2(\omega n n^*) = f_2(\omega p p^*),$$

(5)

where $n^*$ is the $n$-like member of the antidecuplet. The proton-neutron relation for the $\phi$-couplings is similar to that for $\omega$.

In the same manner, one can find $K^*$-couplings for transitions between the nucleon and $\Sigma^*$, the $\Sigma$-like members of the antidecuplet. According to usual isotopic relations

$$f_2(\bar{K}^*0 p \Sigma^{*+}) = -f_2(K^*- n \Sigma^{*-}) = -\sqrt{2} f_2(K^*- p \Sigma^{*0}) = \sqrt{2} f_2(\bar{K}^*0 n \Sigma^{*0}),$$

(6)

while $U$-spin relations give

$$f_2(K^*0 p \Theta^+) = -\sqrt{3} f_2(\bar{K}^*0 p \Sigma^{*+}).$$

(7)

Evidently, all the couplings may indeed be expressed through one of them. We can take, for example, $f_2(\rho^0 n n^*)$ and then find

$$f_2(\bar{K}^*0 p \Sigma^{*+}) = -\sqrt{2} f_2(\rho^0 n n^*), \quad f_2(K^*- p \Sigma^{*0}) = f_2(\rho^0 n n^*).$$

(8)

In turn, the $\rho$-meson coupling may be estimated starting from the transition magnetic moment.

### 4 Vector meson dominance

A good approximation for non-hard electromagnetic interactions of hadrons is the hypothesis that “the entire hadronic electromagnetic current operator is identical with a linear combination of the known neutral vector–meson fields” [32]. This approach is now known as the vector–meson dominance (VMD) [33] (for a recent brief review see the talk [34]).

The simplest form of VMD takes into account only the lightest mesons $\rho^0$, $\omega$, and $\phi$. Then the transition magnetic moments may be expressed as

$$\mu(B_1 \to B_2) = \frac{1}{m_1 + m_2} \sum_{V = \rho^0, \omega, \phi} \frac{e}{g_V} f_2(V B_1 B_2) |_{q^2 = 0}. \quad (9)$$

In what follows, we will use $f_2(V B_1 B_2)$ only at $q^2 = 0$, without showing this explicitly every time.

The meson-photon couplings $g_V$ can be easily related to the partial widths of decays $V \to e^+e^-$:

$$\frac{g_V^2}{4\pi} = \frac{\alpha^2}{3} \frac{m_V}{\Gamma(V \to e^+e^-)}. \quad (10)$$

4
Assuming exact $SU(3)_F$, the photon corresponds to the octet component with $U = 0$. The couplings $g_V$ then satisfy the group relations

$$\frac{1}{g_\rho} : \frac{1}{g_\omega} : \frac{1}{g_\phi} = 1 : \frac{1}{3} : \left(-\frac{\sqrt{2}}{3}\right). \quad (11)$$

For illustration, let us consider how these relations, together with other $SU(3)_F$ relations \(3\) and VMD relation \(5\), result in the cancellation of various contributions to $\mu(p^* \rightarrow p)$:

$$(m_p + m_{p^*}) \cdot \mu(p^* \rightarrow p) = \frac{e}{g_\rho} f_2(\rho^0 p p^*) \left(1 - \frac{1}{3} - \frac{2}{3}\right) = 0. \quad (12)$$

Relations \(11\), with an additional assumption $m_V^{(0)} = m_V^{(8)}$, also predict that

$$\Gamma(\rho^0 \rightarrow e^+ e^-) : \Gamma(\omega \rightarrow e^+ e^-) : \Gamma(\phi \rightarrow e^+ e^-) = 9 : 1 : 2. \quad (13)$$

Experimentally \(2\), these ratios are

$$(7.0 \text{ keV}) : (0.6 \text{ keV}) : (1.3 \text{ keV}) = 11.6 : 1 : 2.1.$$\]

Evidently, $SU(3)_F$-violations are here less than 30%. Taking into account the difference of masses for $\rho^0$, $\omega$, and $\phi$, one can see that the accuracy of ratios Eq. \(11\) is also not worse than 30%. Having in mind the current large uncertainties in properties of the $N^*(1675)$, which we assume to be the $N$-like partner of the $\Theta^+$, we will take the $SU(3)_F$ relations to be exact.

Then, similar to the proton transition moment \(12\), we obtain the neutron transition moment

$$(m_n + m_{n^*}) \cdot \mu(n^* \rightarrow n) = \frac{e}{g_\rho} f_2(\rho^0 p p^*) \left(-1 - \frac{1}{3} - \frac{2}{3}\right) = -\frac{2e}{g_\rho} f_2(\rho^0 p p^*) = \frac{2e}{g_\rho} f_2(p^0 n n^*). \quad (14)$$

### 5 Numerical estimates

We are now ready to discuss the numerical values of various quantities. Experimental characteristics of the $\rho^0$ \(2\) and Eq. \(10\) give $|g_V| \approx 5$. Previously, in Ref. \(23\), we extracted $|\mu(n^* \rightarrow n)| = (0.13-0.37) \mu_N$, where $\mu_N = e/(2m_N)$ is the standard nuclear magneton. This gives

$$|f_2(\rho^0 n n^*)| = |f_2(\rho^0 p p^*)| \approx (0.13 - 0.37) g_\rho m_N + m_{N^*} \approx (0.45-1.28), \quad (15)$$

which is essentially smaller than $|f_2(\rho^0 N N)|$, equal to $12 - 16$ \(28\).

With the value in Eq. \(15\), Eqs. \(1\) and \(5\) give

$$|f_2(K^{*0} p \Theta^+)| = |f_2(K^{*+} n \Theta^+)| = \sqrt{6} |f_2(\rho^0 n n^*)| = (1.10 - 3.14). \quad (16)$$

Analogously, from Eqs. \(13\) and \(15\), we can find $K^*$-couplings for the $\Sigma^*$-excitation:

$$|f_2(K^{*-} p \Sigma^{0})| = (0.45 - 1.28), \quad |f_2(K^{*0} p \Sigma^{*0})| = (0.64 - 1.81). \quad (17)$$

\(^1\)In exact $SU(3)_F$, $\mu(p^* \rightarrow p)$ should vanish, since the $U$-spins are $3/2$ for $p^*$, $1/2$ for $p$, and $0$ for the photon. With the violation of $SU(3)_F$, this transition moment becomes non-vanishing, but is still much smaller than $\mu(n^* \rightarrow n)$ \(33\).
Incidentally, the value of $1.1$ for $|f_2(K^+n\Theta^+)| = |f_2(K^0p\Theta^+)|$, supported now by Eq. (16), was used earlier in Ref. [22] to estimate the production cross section of $\Theta^+$ in photoreactions. With a $\Theta^+$-width of $1$ MeV (in accordance to Refs. [36]) and the above value of $f_2(K^0N\Theta)$, calculations [22] find small cross sections $\sigma_{tot}(\gamma p \rightarrow K^0\Theta^+) < 0.22 \text{ nb}$ and $\sigma_{tot}(\gamma n \rightarrow K^-\Theta^+) < 1 \text{ nb}$, which are below the limits given recently by the CLAS Collaboration [4, 6].

Note the difference between proton and neutron targets. Photoproduction off the proton is nearly independent of the width of $\Theta^+$, due to absence of the $K^0$-exchange. Therefore, the cross section of $\gamma p \rightarrow K^0\Theta^+$ is mainly determined by the $K^0$-exchange, which is proportional to $|f_2(K^0p\Theta^+)|^2$ (for more details, see Ref. [22]). On the other side, the neutron photoproduction $\gamma n \rightarrow K^-\Theta^+$ is mainly determined by the charged $K$-exchange, which is proportional to $\Gamma_{\Theta^+}$.

To clarify the meaning of our numerical values for $f_2$, we consider in more detail the photoproduction

$$\gamma + p \rightarrow K^0 + \Theta^+, \quad (18)$$

with $K^*$-exchange as the main contribution. We can compare experimental limits for this reaction, obtained by the CLAS Collaboration [6], and the theoretical calculations in the model of Ref. [22]. Fig. 1 shows the total cross section of photoproduction (18) as a function of the photon energy. Both experimental and theoretical results shown here assume the same value $m_{\Theta^+} = 1540$ MeV. The filled circles correspond to the experimental upper limits [6] for particular initial energies. The lower edge of the shaded area reproduces the lower curve in Fig.7 of Ref. [22], for $(J^P)_{\Theta^+} = (1/2)^+$. To obtain the whole shaded area in our Fig. 1 we multiply this curve by the factor $|f_2(K^*p\Theta^+)|/1.1|^2$, with $f_2(K^*p\Theta^+)$ running the whole above range in Eq. (16) for the $K^*$ coupling constant.

We see that, apart from the higher photon energy region and/or the upper edge of the interval Eq. (16) for $f_2$, the estimate of the cross section is below the limits put by the CLAS Collaboration. The theoretical results are very rough, of course, and contain various uncertainties, both in the approximations used and in the phenomenological inputs. Nevertheless, the comparison in Fig. 1 shows that the CLAS analysis procedure of Ref. [6] might not be sensitive enough to reveal the $\Theta^+$ (if it exists).

Futhermore, Fig. 1 clearly shows that the $K^*$ coupling to the $N\Theta^+$ system should be very small indeed. Moreover, we expect that corrections due to a nonzero value of the strange quark mass may further reduce our present estimates.

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Figure 1: Shaded band corresponds to the range of $\Theta^+$ production cross section in the process $\gamma + p \rightarrow K^0 + \Theta^+$ as a function of the photon energy, expected according to model [22]. The lower and upper edges of the band correspond to the coupling constant values $|f_2(K^*0\, p\, \Theta^+)| = 1.10$ and $3.14$, the range obtained in the present paper, Eq. (16). The filled circles correspond to the experimental upper limits for the cross section given by the CLAS Collaboration [6].

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