Anomalous Hypercharge Axial Current And The Couplings Of
The $\eta$ And $f_1(1420)$ Mesons To The Nucleon

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Abstract

It is emphasized that the precise three flavor symmetry of hadrons is not SU(3)$_F$ but rather U(4)$_F$ restricted to SU(2)$_{ud}$⊗SU(2)$_{cs}$⊗U(1) and considered in the limit of frozen charm degree of freedom. Within this scheme the hypercharge generator is necessarily an element of the su(2)$_{cs}$⊗u(1) subalgebra as it contains the baryon number generator associated with the unit matrix. The structure of hypercharge obtained in this way is the only one that is consistent with the Gell-Mann–Nishijima relation. In considering now the corresponding axial hypercharge transformations, the unit element of the u(4) algebra will give rise to the anomalous U(1)$_A$ current and the resulting hypercharge axial current will be anomalous, too. It is shown that the only anomaly free neutral strong axial current having a well defined chiral limit is identical (up to a constant factor) with the weak axial current. There, a purely strange axial current comes in place of the hypercharge one and the $\eta$ meson acquires features of a ‘masked’ strange Goldstone boson. The consequence is that the $\eta N$ and $f_1(1420)N$ couplings have to proceed via a purely strange isosinglet axial current. Therefore, the $\eta$ and $f_1(1420)$ mesons probe at the tree level the polarization of the strange quark sea and their couplings
appear strongly suppressed relative to quark model predictions. For this rea-
son, loop vertex corrections acquire importance. A model based on effective 
lagrangians for coupling the $f_1(1420)$ and $\eta$ mesons to the nucleon via trian-
gular vertex corrections containing two–meson states has been developed and 
shown to be convenient for data description beyond the limits of applicability 
of chiral perturbation theory.

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and $f_1$ meson-nucleon couplings
I. INTRODUCTION

Three flavor SU(3)$_F \otimes$SU(3)$_F$ chiral symmetry (ChS) of strong interaction is one of the basic guidelines in constructing the nucleon dynamics as it is considered to be the global internal symmetry of the QCD lagrangian [1]. As long as ChS is supposed to be realized in the non–multiplet Nambu–Goldstone mode with the pseudoscalar octet mesons acting as the associated Goldstone bosons, it becomes possible to expand correlation functions in powers of the current quark masses and the external Goldstone boson momenta thought to be small at the hadronic scale of $\Lambda \sim 1$ GeV [2]. This so–called chiral perturbation theory (ChPT) advanced in the last decade to a powerful scheme for describing various low–energy phenomena [3]. In view of this success, establishing the Goldstone boson character of the lowest pseudoscalar mesons attracts special attention.

The first aim of the present study is to emphasize that the $\eta$ meson cannot be considered as the flavor octet Goldstone boson but is rather a ‘would be’ strange Goldstone boson, a result already conjectured in the previous work [4]. The second is to suggest a phenomenological field theoretical approach for describing meson–nucleon vertices beyond the limits of applicability of ChPT. The reason for the non–octet Goldstone boson nature of the $\eta$ meson is that according to an observation reported in Ref. [5], the precise three flavor symmetry of hadrons is not SU(3)$_F$ but rather U(4)$_F$ restricted to SU(2)$_{ud} \otimes$SU(2)$_{cs} \otimes$U(1) and considered in the limit of frozen charm degree of freedom. In such a case, the only anomaly free neutral strong axial current of a well defined chiral limit is obtained in forbidding the U(1)$_A$ axial transformation. This current has same flavor structure as the neutral weak axial current. Thus, within the three flavor space, a purely strange isosinglet axial current comes in place of the flavor octet axial current, and no hypercharge Goldstone boson is required any more. Rather, a purely strange Goldstone boson should be introduced, whose appearance is however prevented through the violation of the OZI rule for the pseudoscalar mesons as brought about by the U(1)$_A$ anomaly.

The consequence is that the neutral axial and pseudoscalar mesons will couple to the
nucleon via the \( \bar{s}s \) quarkonium components of their wave functions. The corresponding point-like vertices will be proportional to \( \Delta s \), the small fraction of proton helicity carried by the strange quark sea, and appear strongly suppressed relative quark model predictions. In view of that smallness, loop vertex corrections acquire importance.

The effect of the one-loop corrections on the experimentally observed strong suppression of the nucleon matrix element of the octet (or, hypercharge) axial vector current has been considered, for example, in Ref. [6] within the framework of SU(6) chiral perturbation theory. There, the strong dependence of the result on the \( N-\Delta \) mass splitting was revealed and the necessity for higher-order corrections discussed. Instead of considering the suppression of the nucleon octet axial matrix element, we here rather examine the small enhancement of the axial \( (\bar{s}s)N \) coupling in terms of triangular corrections of the type \( a_0(980)\pi N \) to the \( \eta NN^- \), and \( K^*(892)KY \) to the \( f_1(1420)NN \) vertices, respectively. Such triangular vertices participate effective \( Z(\bar{s}s)N \) chain couplings. Indeed, the coupling of an external weak isosinglet axial current to the nucleon can be viewed as being mediated by the \( f_1(1420) \) or \( \eta \) isosinglet meson states which contain strange and non-strange quarkonia simultaneously, thus violating the OZI rule. In such a case, the \( Z \) boson has at the weak vertex the opportunity to select the strange \( \bar{s}s \) quarkonium from the wave function of the respective intermediate meson, while at the strong vertex the nucleon can couple to the remaining \( (\bar{u}u+\bar{d}d) \) quarkonia via its meson cloud as parametrized by the triangular vertices introduced above.

A method similar in spirit to the one presented here but of different techniques is the so-called meson cloud model of Ref. [7]. There, the authors study the influence of the mesons surrounding the nucleon on its electromagnetic and weak vector form factors in terms of a nucleon wave function having a non-negligible overlap with the nucleon-meson scattering continuum. We here focus rather on the isoscalar axial form factor.

The paper is organized as follows. In the next section we consider the identification criteria for the Goldstone boson character of the \( \eta \) meson, discuss the ambiguities related to the calculation of the \( \eta \) meson pole term, and show that the \textit{axial hypercharge current is anomalous}. In Sect. 3 the contact \( \eta NN \) and \( f_1(1420)NN \) vertices are considered. In Sect.
4 the effective $\eta NN$ vertex is calculated from the $\pi a_0(980)N$ triangular vertex, while Sect. 5 contains the results on the effective $f_1(1420)$ meson nucleon-couplings associated with the $KK^*(892)Y$ triangle. The paper ends with a short summary.

II. $\eta$ GOLDSTONE BOSON REVISITED

The $\eta$ meson is canonically considered as the octet Goldstone boson of three flavor $SU(3)_L \otimes SU(3)_R$ chiral symmetry as the sum of its pole term current, displayed in Fig. 1, and the hypercharge axial current of the nucleon ($N$) is assumed to be partially conserved [8]. The basic tacit assumption entering the calculation of the $\eta$ pole term is the universality of the hypercharge axial current at both the strong and weak vertices. Let us consider as an illustration the canonical ansatz according to which both the neutral flavor axial current and the $\eta$ meson wave function are approximately determined by Gell-Mann’s hypercharge matrix $\lambda^8$, i.e.

$$
|\eta\rangle \approx \frac{\lambda^8}{\sqrt{2}} q_3, \quad j_{\mu,5}^\eta = \bar{q}_3 \gamma_\mu \gamma_5 \frac{\lambda^8}{2} q_3, \quad \lambda^8 = \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix}.
$$

We denote by $q_3$ the quark field in three flavor space: $q_3 = (u d s)^T$. The $\eta$ weak decay current is now parametrized as

$$
J_{\mu,5}^\eta := \langle 0 | j_{\mu,5}^\eta | \eta \rangle = f_\eta q_\mu.
$$

Here, $q_\mu$ and $f_\eta$ stand for the four momentum and weak decay coupling constant of the $\eta$ meson, respectively. In inserting Eq. (1) into (2) and assuming in accordance to Ref. [9] the quarkonia-quark currents couplings to be diagonal in flavor, i.e.

$$
\langle 0 | \frac{1}{2} \bar{q} q | \bar{q}' q' \rangle = \delta_{qq'} \kappa_q (0^-), \quad q, q' = u, d, s,
$$

the value of $f_\eta$ is calculated as

$$
3\sqrt{2} f_\eta = \kappa_u(0^-) + \kappa_d(0^-) + 4\kappa_s(0^-).
$$
Within this scheme, the contact octet meson–nucleon vertex $\mathcal{V}_{\eta NN}^8$ is given the form

$$\mathcal{V}_{\eta NN}^8 = \frac{1}{f_{\eta}^2} J_5^8 \cdot J_5^\eta = \frac{1}{f_{\eta}^2} \left( G_A^8 \bar{N}_2 \gamma_5 \frac{\mathbb{I}}{2} N_1 \right) \cdot (f_{\eta} i q \phi_{\eta}) \ .$$

(5)

Here, $\phi_{\eta}$ stands for the field of the $\eta$ meson, while $N_1$ and $N_2$ denote the spinor fields of the in– and outcoming nucleons, respectively. The parameter of the dimension $[mass^2]$ entering the contact current–current coupling equals the weak $\eta$ meson decay coupling constant. This choice implies the *current universality* mentioned above, as it allows one to express the coupling of the $\eta$ meson to arbitrary baryon targets in terms of its weak decay coupling $f_{\eta}$.

The quantity $G_A^8$ in (5) stands for the coupling of an external axial lepton current to the nucleon hypercharge axial current $J_{\mu,5}$ and is defined via

$$s^\mu J_{\mu,5}^\eta := s^\mu \langle N_2 | J_{\mu,5}^8 | N_1 \rangle = G_A^8 \bar{U}_{N_2} \gamma_\mu \gamma_5 \frac{\mathbb{I}}{2} U_{N_1} s^\mu, \quad s^\mu q_\mu = 0 ,$$

(6)

with $s^\mu$ being a unit spin polarization vector and $U_N$ denoting nucleonic spinors. The value of the nucleon hypercharge coupling $G_A^8$ following from Eq. (6) is expressed in terms of the helicity fractions $\Delta u$, $\Delta d$, and $\Delta s$, carried by the respective $u$, $d$, and $s$ quarks as

$$G_A^8 = \frac{1}{\sqrt{3}} (\Delta u + \Delta d - 2\Delta s) .$$

(7)

The combination $\frac{G_A^8}{2f_{\eta}}$ in Eq. (6) is conventionally denoted by $\frac{f_{NN\eta}}{m_{\eta}}$ with $f_{NN\eta}$ being called the gradient (or pseudovector (PV)) $\eta N$ coupling constant i.e.

$$\frac{f_{NN\eta}}{m_{\eta}} = \frac{G_A^8}{2f_{\eta}} = \sqrt{\frac{3}{2}} \frac{\Delta u + \Delta d - 2\Delta s}{\kappa_u(0^-) + \kappa_d(0^-) + 4\kappa_s(0^-)} ,$$

(8)

where use has been made of Eqs. (4) and (7). On-mass shell, the PV coupling can be expressed through the pseudoscalar (PS) coupling $g_{\eta NN}$ via the equivalence relation [8] as

$$\frac{f_{NN\eta}}{m_{\eta}} = \frac{g_{NN\eta}}{2m_N} = \frac{G_A^8}{2f_{\eta}} ,$$

(9)

with $m_N$ standing for the nucleon mass. The constraint on $g_{NN\eta}$ in Eq. (8) is known as the hypercharge ‘Goldberger-Treiman’ (GT) relation. Ordinarily, the GT constraint is alternatively obtained rather as a condition sufficient for the partial conservation of the
hypercharge axial current already at the tree level and is indicative for the Goldstone-boson character of the $\eta$ meson. The considerations given so far illustrate that it is, actually, the assumed universality of the hypercharge axial current at both the strong and weak vertices of the $\eta$ pole term which underlies the hypercharge GT-relation and thus enables the realization of hypercharge chiral symmetry in the hidden Nambu–Goldstone mode, and vice versa. In noting now that the weak neutral axial current $J_{\mu,5}^w$,

$$J_{\mu,5}^w = -\frac{1}{2}(\bar{u}d)\gamma_\mu\gamma_5\frac{\tau_3}{2}\begin{pmatrix} u \\ d \end{pmatrix} + \frac{1}{4}\bar{s}\gamma_\mu\gamma_5 s.$$ (10)

apparently differs from $J_{\mu,5}^8$, one immediately realizes that the universality assumption brings an element of ambiguity in the calculation of the $\eta$ pole term. Indeed, in assuming the $\eta$ meson to couple to the hadronic vacuum via the hypercharge axial current (cf. Eq. (2)), one ascribes to this meson the ability to decompose the fundamental Z-boson current into hypercharge and flavor singlet components, much like an optically active material decomposes a linearly polarized wave into its circularly polarized parts. In that sense the neutral axial and pseudoscalar mesons are considered as a sort of ‘electroweak active’ probes. This means that a constituent symmetry is given higher priority over a gauge symmetry without any deeper justification. Remarkably, the assumed priority does not find any confirmation by data which speak in favor of a surprisingly small $\eta N$ coupling rather than in favor of the corresponding GT-prediction (see [4] for details) even after accounting for the small $\eta - \eta'$ mixing required by the Gell-Mann-Okubo mass formulae. That such a mixing will lead to a reduction of the nucleon matrix element of the octet axial current was earlier considered, among others, in Ref. [10]. There, however, the reduction reported was far away from being as big as the required one. Note, however, that the accuracy of GT-relations has been reliably proven only for the case of the pion–nucleon system [11]. Below we argue that it is actually the electroweak axial current that enters the calculation of pole terms containing neutral pseudoscalar and axial vector mesons. The argumentation given below essentially follows Ref. [4].
Consider the fundamental four flavor vector current of the quarks

\[ j_\mu = \frac{2}{3} \bar{u}_\mu u - \frac{1}{3} \bar{d}_\mu d + \frac{2}{3} \bar{c}_\mu c - \frac{1}{3} \bar{s}_\mu s \]

\[ = \bar{q} t_3 \gamma_\mu q + \bar{q} \frac{Y}{2} \gamma_\mu q, \quad q = (u \, d \, c \, s)^T. \]  

(11)

Here, \( Y \) and \( t_3 \) in turn represent hypercharge and third projection of isospin within the quark flavor quadruplet, \( q \), and are explicitly given below as

\[ t_3 = \frac{1}{2} \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}, \quad Y = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & \hat{C} & 0 \\ 0 & 0 & \hat{S} & 0 \end{pmatrix} + \frac{1}{3} \mathbb{I}_4, \]  

(12)

with \( \hat{C}c = c \), and \( \hat{S}s = -s \), respectively. As long as \( j_\mu \) is conserved, its total charge

\[ Q(t) = \int j_0(t, \vec{x}) d^3 \vec{x}, \]  

(13)

is a constant of motion and labels the hadron states. When considered as an operator, \( \hat{Q} \) is directly read off from Eq. (11) to be related to the operators of isospin \( \hat{t}_3 \) and hypercharge \( \hat{Y} \) via the famous Gell-Mann–Nishijima relation

\[ \hat{Q} = \hat{t}_3 + \frac{1}{2} \hat{Y}. \]  

(14)

Now if Gell-Mann’s matrix \( \lambda^8 \) in Eq. (11) were to be interpreted as the generator of three flavor hypercharge, it should emerge in the limit of negligible \( c \) quark effects from the complete four flavor hypercharge \( Y \) of Eq. (12) as

\[ \frac{1}{3} \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & -2 \end{pmatrix} = \lim_{m_c \to \Lambda_c} \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & \hat{C} & 0 \\ 0 & 0 & \hat{S} & 0 \end{pmatrix} + \frac{1}{3} \lim_{m_c \to \Lambda_c} \mathbb{I}_4. \]  

(15)

Here \( \Lambda_c \) has to be sufficiently large in order to ‘freeze out’ the charm degree of freedom on the 1 GeV mass scale, on the one side, and still finite, in order to preserve the anomaly
free character of the SU(4)$_F$ theory [12], on the other side. From the latter equation one sees that the matrix $\lambda^8$ representing the hypercharge in the truncated flavor space (now three dimensional) happens by accident to be traceless while the full four flavor hypercharge matrix $Y$ in Eq. (12) is not traceless as it contains the unit matrix corresponding to the baryon number current. In this way the misleading impression appears that hypercharge can be introduced on the level of the group SU(3)$_F$ and be exploited for the construction of a partially conserved axial current. In contrast to Eq. (15), Gell-Mann’s hypercharge matrix $\lambda^8$ from the su(3) algebra decomposes into

$$\sqrt{3} \lambda^8 = \lambda^3 + 2 \lambda_U^3,$$

$$\lambda^3 = \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad \lambda_U^3 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -1 \end{pmatrix}, \quad (16)$$

with $\frac{1}{2} \lambda^3$ and $\frac{1}{2} \lambda_U^3$ in turn denoting third projections of isospin and $U$-spin$^9$. Looking at Eq. (16), one finds the following expression for the electric charge of the quarks,

$$Q = t_3 + \frac{1}{3} (t_3 + 2t_U^3), \quad (17)$$

with

$$t_U^3 = \frac{1}{2} \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix}. \quad (18)$$

Eq. (17) differs from the standard relation, $Q = t_3 + (S + B)/2$, in group theoretical aspects. It may be recalled that the full Gell-Mann–Nishijima relation, $Q = t_3 + (S + C + B)/2$ is a genuine $u(1)_B \oplus su(4)_F$-relation with the sum $S + C + B$ being the four flavor hypercharge $Y$.

$^1$The matrices $\lambda^3$ and $\lambda_U^3$ are nothing but the diagonal elements of the su(3) algebra in the so-called Weyl basis.
In the limit of neglected $c$ quark effects, this relation reduces to the form $Q = t_3 + (S + B)/2$ which, however, cannot be read as an $\text{su}(3)_F$-relation since the baryon number $B$ becomes extraneous to this context. Thus, $(S + B)/2$ does not correspond to any $\text{su}(3)_F$ generator, none of the $\text{su}(3)_F$ algebra elements can be given the interpretation of a hypercharge generator in the limit case of three flavors. The physical hypercharge generator appears as a non–traceless element of the vector space spanned by the $\text{u}(4)_F$ algebra and the corresponding hypercharge axial current is anomalous. Indeed, in considering axial hypercharge transformations, the term containing the unit matrix on the rhs of Eq. (15) will give rise to the anomalously divergent $\text{U}(1)_A$ current for which no chiral limit can be formulated [13].

As a consequence, the hypercharge axial current will be anomalous, too. The only neutral flavor axial current which appears to be conserved in the chiral limit of vanishing quark masses will be $j_{\mu,5}$ defined as

$$j_{\mu,5} = \bar{q} \gamma_\mu \gamma_5 \left( t_3 + \frac{1}{2} (Y - \frac{1}{3}) \right) q . \tag{19}$$

The flavor structure of $j_{\mu,5}$ in the last equation reflects the exclusion of the anomalous $\text{U}(1)_A$ current [13] which cannot be used any longer as a building block for the construction of an anomaly free octet axial current. One remarkable feature of $j_{\mu,5}$ is that its structure is identical (up to the factor of $-1/2$) to that of the neutral weak axial vector current and respects the OZI rule. For this reason, the well established universality of the flavor changing weak and strong axial vector currents underlying the current algebra can be extended to include the neutral ones. It is this current which will enter the calculations of pole terms created by neutral pseudoscalar and axial vector mesons, and the above mentioned ambiguity is resolved by now. The $j_{\mu,5}$ current decomposes in three flavor space into an isovector ($j^I_{\mu,5}$) and a purely strange $\text{SU}(2)_I$ isosinglet ($j^s_{\mu,5}$) component

$$j_{\mu,5} = j^I_{\mu,5} + j^s_{\mu,5} ,$$

$$j^I_{\mu,5} = \bar{q}_3 \gamma_\mu q_3 , \quad j^s_{\mu,5} = -s \gamma_\mu \gamma_5 \frac{1}{2} s , \tag{20}$$

with $\lambda^3$ being the isospin Gell-Mann matrix $\lambda^3 = \text{diag}(1,-1,0)$. 

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A remark is worthy on the flavor structure of the \( \eta \) meson which is the pseudoscalar analogue to the vector meson \( \phi \). While \( \phi \) is an almost pure strange quarkonium, the wave function of the physical \( \eta \) meson derived from fitting the meson mass spectrum has a significant non-strange quarkonium component according to

\[
|\eta\rangle = \cos \epsilon (-|\bar{s}s\rangle) - \sin \epsilon \frac{1}{\sqrt{2}}(|\bar{u}u + \bar{d}d\rangle), \quad \epsilon = -45.4^\circ. \tag{21}
\]

Here, \( \frac{1}{\sqrt{2}}(\bar{u}u + \bar{d}d) \) is the singlet \( U(2)_I \) state. For \( \epsilon = -\arctan 1/\sqrt{2} \) the octet scalar state \( |\eta_8\rangle = q_3 \frac{A^s}{\sqrt{2}} q_3 \) is reproduced. Eq. (21) clearly illustrates that the wave function of the \( \eta \) meson transforms in accordance with a representation of \( \lim_{m_c \to \Lambda_c} U(1)_I \otimes SU(2)_{cs} \) rather than as a genuine \( SU(3)_F \) state. From this point of view the predominantly scalar octet nature of the \( \eta \) meson appears as an artefact of the violation of the OZI rule for the pseudoscalar mesons as brought about by the \( U(1)_A \) anomaly [14] rather than through a fundamental underlying \( SU(3)_F \) symmetry.

Correspondingly, the \( \eta' \) function reads,

\[
|\eta'\rangle = \sin \epsilon (-|\bar{s}s\rangle) + \cos \epsilon \frac{1}{\sqrt{2}}(|\bar{u}u + \bar{d}d\rangle), \quad \epsilon = -45.4^\circ. \tag{22}
\]

The large value of the angle \( \epsilon \) in the last two equations signals a much stronger violation of the OZI rule within the pseudoscalar nonet as compared to the vector meson nonet, where \( \epsilon \approx 5^\circ \).

III. CONTACT \( \eta NN \) AND \( F_1(1420)NN \) VERTICES

The absence of a hypercharge component in the anomaly free neutral strong axial current does not contradict the fact that the structure of the \( \eta \) meson as deduced from data fits by means of the Gell-Mann-Okubo mass formulae, deviates from the purely strange quarkonium. It only has essential impact on the flavor structure of the vertices including neutral octet axial and pseudoscalar mesons such as the \( \eta \) and \( f_1(1420) \) mesons to the nucleon. In other words, while the leading components of the \( \eta \) and \( f_1(1420) \) meson wave functions can
still be the scalar and singlet octet states, respectively, the $\eta NN$- and $f_1(1420)NN$ vertices will rather be purely strange isosinglets than of $F$-type because both these mesons can couple only to the anomaly free strange isosinglet nucleon current given below as

$$J_{\mu,5}^{s(N)} = \langle N | \frac{1}{2} \bar{s} \gamma_\mu \gamma_5 s | N \rangle = G_1^s \bar{U}_N \gamma_\mu \gamma_5 \frac{1}{2} U_N , \quad G_1^s = \Delta s ,$$

which can happen only via their strange quarkonium ingredients.

In approximating, for simplicity, the flavor part of the wave functions of both the $\eta$ and $f_1(1420)$ mesons by Eq. (21) without caring at the moment about the concrete numerical value of the corresponding mixing angle, their new axial currents are now defined by

$$J_{\mu,5}^{s(\eta)} = \langle 0 | - \frac{1}{2} \bar{s} \gamma_\mu \gamma_5 s | \eta \rangle = f_\eta i q_\mu , \quad f_\eta := \cos \epsilon \kappa_s(0^-) m_\eta ,$$

$$J_{\mu,5}^{s(f_1)} = \langle 0 | - \frac{1}{2} \bar{s} \gamma_\mu \gamma_5 s | f_1 \rangle = f_{f_1} m_{f_1}^2 \epsilon_{f_1,\mu} , \quad f_{f_1} := \cos \epsilon \kappa_s(1^+).$$

Here $\kappa_s(0^-)$ and $\kappa_s(1^+)$ denote the dimensionless couplings of the strange quarkonium components of the pseudoscalar and axial vector mesons to the strange axial vector current, respectively. The quantities $f_{f_1}$, $m_{f_1}$, and $\epsilon_{f_1,\mu}$ stand for the weak decay coupling, the mass and the polarization vector of the $f_1(1420)$ meson. Furthermore, the expressions in Eq. (24) have been obtained in assuming universal couplings of the $\bar{u}u$ and $\bar{d}d$ quarkonia to the quark axial currents, i.e. $\kappa_u(0^-) = \kappa_d(0^-)$, and $\kappa_u(1^+) = \kappa_d(1^+)$, to recover isospin symmetry already at the tree level. In accordance with Eq. (5) the purely strange isosinglet contact $\eta NN$ and $f_1(1420)NN$ vertices now read

$$V_{\eta NN}^s = \frac{1}{f_\eta^2} J_{s(\eta)}^s \cdot J_{s(\eta)}^s = \frac{1}{f_\eta^2} G_1^s \bar{N} \gamma \gamma_5 \frac{1}{2} N \cdot f_\eta i q_\eta \phi_\eta ,$$

$$V_{f_1 NN}^s = \frac{1}{m_{f_1}^2} J_{s(f_1)}^s \cdot J_{s(f_1)}^s = \frac{1}{m_{f_1}^2} G_1^s \bar{N} \gamma \gamma_5 \frac{1}{2} N \cdot f_{f_1} m_{f_1}^2 \phi_{f_1} ,$$

with $\phi_{f_1}$ standing for the $f_1$ meson field. The combination $\frac{G_1^s f_{f_1}}{2}$ is ordinarily identified with the $f_1(1420)N$ contact coupling $f_{f_1 NN}$. In inserting Eqs. (23) and (24) in the last expressions, one is led to the following relations:

$$2 \text{ The empirically observed closeness of the } \pi \text{ and } \eta \text{ weak decay constants, } (f_\eta \approx 1.1 f_\pi) \text{ does not necessarily imply } \kappa_u(0^-) \approx \kappa_s(0^-) \text{ at the tree level.}$$
\[ f_{\eta NN} = \frac{G_s^s}{2 f_{\eta}}, \quad \text{with} \quad f_{\eta NN} = \frac{\Delta s}{2 \cos \epsilon \kappa_s(0^-)}, \]
\[ f_{f_1 NN} = \frac{G_s^s f_1}{2}, \quad \text{with} \quad f_{f_1 NN} = \frac{\Delta s}{2 \cos \epsilon \kappa_s(1^+)}. \quad (26) \]

In contrast to Eq. (7), the polarization of the non–strange sea does not any longer participate the weak decay couplings of the \( \eta \) and \( f_1(1420) \) mesons. In exploiting the on–shell equivalence relation between pseudoscalar and pseudovector couplings \[8\] one finds for the pseudoscalar \( \eta N \) coupling constant \( g_{\eta NN} \) the following expression
\[ g_{\eta NN} = \frac{f_{\eta NN}}{m_\eta} 2m_N = \frac{\Delta s}{2 \cos \epsilon \kappa_s(0^-)} \frac{2m_N}{m_\eta}. \quad (27) \]

As long as Eq. (26) links the \( \eta N \) coupling to \( G_s^s \) by means of a Goldberger-Treiman relation, the \( \eta \) meson can be considered as a ‘would be’ strange Goldstone boson. This means that at tree level the contact \( \eta N \) coupling appears proportional to the fraction of proton helicity carried by the strange quark sea, rather than to the octet axial vector coupling \( G_A^{(8)} = \frac{1}{\sqrt{3}}(\Delta u + \Delta d - 2\Delta s) \), a result already conjectured in a previous work \[4\]. Due to the smallness of \( \Delta s = -0.08 \pm 0.05 \) (see \[15\] for a recent review), the \( \eta \) and \( f_1(1420) \) mesons will almost decouple at tree level from the nucleon. This might be one of the main reasons for which a strong suppression of the \( \eta N \) couplings has frequently been found over the years by various data analyses of \( \eta \) photoproduction off proton at threshold \[16\], \( \bar{p}p \) collisions \[17\], as well as nucleon-nucleon (NN) and nucleon–hyperon (NY) phase shifts \[18\].

One of the main points of the present study is that the tree level \( \eta N \) and \( f_1 N \) couplings almost vanish as they proceed over a purely strange isosinglet axial current. In the following, the small but non–negligible couplings of these mesons to the nucleon will be entirely attributed to triangular vertices of the type \( a_0 \pi N \), and \( KK^*Y \), respectively.

IV. EFFECTIVE \( \eta NN \) VERTICES

In this section we calculate the gradient and pseudoscalar coupling constants of the \( \eta \) meson to the nucleon by means of triangular vertices involving the \( a_0(980) \) and \( \pi \) mesons.
The special role of the $a_0(980)\pi N$ triangular diagram as the dominant one-loop mechanism for the $\eta N$ coupling is singled out by the circumstance that the $a_0(980)$ meson is the lightest meson with a two particle decay channel containing the $\eta$ particle [19]. The contributions of heavier mesons such as the isotriplet $a_2(1320)$ tensor meson with an $\eta\pi$ decay channel and the isoscalar $f_0(1400)$, $f'_2(1525)$ and $f_2(1720)$ tensor mesons with $\eta\eta$ decay channels will be left out of consideration because of the short range character of the corresponding triangle diagrams on the one side, and because of the comparatively small couplings of the tensor mesons to the nucleon [20,21] on the other side.

The $\pi a_0 N$ triangular couplings (Fig. 2) have been calculated using the following effective lagrangians of common use:

$$\mathcal{L}_{a_0\eta\pi}(x) = f_{a_0\eta\pi} \frac{m_{a_0}^2 - m_\eta^2}{m_\pi} \phi_\eta^\dagger(x) \vec{\phi}_\pi(x) \cdot \vec{\phi}_{a_0}(x)$$

$$\mathcal{L}_{\pi NN}(x) = \frac{f_{\pi NN}}{m_\pi} \bar{N}(x) \gamma_\mu \gamma_5 \vec{\tau} N(x) \cdot \partial^\mu \vec{\phi}_\pi(x),$$

$$\mathcal{L}_{a_0 NN}(x) = g_{a_0 NN} i \bar{N}(x) \vec{\tau} N(x) \cdot \vec{\phi}_{a_0}(x).$$

Here $f_{\pi NN}$ and $g_{a_0 NN}$ in turn denote the pseudovector $\pi N$ and the scalar $a_0 N$ coupling constants. We adopt for $f_{\pi NN}$ the standard value $f_{\pi NN}^2/4\pi = 0.075$ and fit $g_{a_0 NN}$ to data. The value of $f_{a_0\eta\pi} = 0.44$ has been extracted from the experimental decay width [19] when ascribing the total $a_0$ width to the $a_0 \to \eta + \pi$ decay channel. The amplitude $T_{\eta N}(\pi + N \to a_0 + N)$ entering the diagram in Fig. 2 can be parametrized in terms of the following complete set of invariants

$$T_{\eta N}(\pi + N \to a_0 + N) = \bar{U}_N(\vec{p}') \left( \frac{G_1(k^2)}{m_\eta} k_5 \gamma_5 + G_2(k^2) \gamma_5 \right) U_N(\vec{p}) \right),$$

where the invariant functions $G_1(k^2)$ and $G_2(k^2)$ in turn correspond to pseudovector (PV) and pseudoscalar (PS) types of the $\eta N$ coupling. Expressions for $G_1(k^2)$ and $G_2(k^2)$ can be found in evaluating the diagrams in Fig. 3 in accordance with the standard Feynman rules.

3The same argument applies to the neglect of the $f_0(1590)\eta N$ triangular vertex.
We here systematically consider the incoming proton to be on its mass shell and make use of the Dirac equation, so that

$$\gamma_5 \not{p} U_N(\vec{p}) = m_N \gamma_5 U_N(\vec{p}) ,$$

(32) holds. On the contrary, the outgoing proton has been considered to be off its mass shell with

$$\gamma_5 \not{p}' U_N(\vec{p}) = \gamma_5 (\not{p} + \not{k}) U_N(\vec{p}) = m_N \gamma_5 U_N(\vec{p}) + \gamma_5 \not{k} U_N(\vec{p}) .$$

(33)

The final result on $G_1(k^2)$ obtained in this way reads:

$$G_1(k^2) = C \int_0^1 \int_0^1 dy dx \frac{c_1(x, y, k^2)}{Z(m_N, m_\pi, m_{a_0}, x, y, k^2)} ,$$

$$c_1(x, y, k^2) = -\frac{1}{2} x (1 - y) m_N m_\eta ,$$

$$C = \frac{3}{8\pi^2} \frac{m_{a_0}^2 - m_\eta^2}{m_\pi^2} f_{\pi NN} f_{a_0\eta\pi} g_{a_0NN} .$$

(34)

The corresponding expression for $G_2(k^2)$ reads

$$G_2(k^2) = C \int_0^1 \int_0^1 dy dx \frac{c_2(x, y, k^2)}{Z(m_N, m_\pi, m_{a_0}, x, y, k^2)} ,$$

$$c_2(x, y, k^2) = -x (1 - y) m_N^2 .$$

(35)

The function $Z(m_B, m_1, m_2, x, y, k^2)$ appearing in the last two expressions is defined as

$$Z(m_B, m_1, m_2, x, y, k^2) = m_B^2 x^2 (1 - y)^2 + x^2 y k^2 + m_1^2 (1 - x) + (m_2^2 - k^2) xy$$

$$+ (m_B^2 - m_N^2) x (1 - y) .$$

(36)

The remarkable feature of the analytical expressions for the pseudoscalar and pseudovector $\eta N$ couplings is that they are given by completely convergent integrals and depend only on the $a_0 \to \pi + \eta$ decay constant and the respective pion and $a_0$ meson-nucleon couplings.

The sources of uncertainty in the parametrization of the effective $\eta NN$ vertex by means of the triangular $a_0(980)\pi N$ diagram are associated with the $a_0(980)N$ coupling constant and the $\Gamma(\eta\pi)/\Gamma_{a_0}^{tot}$ branching ratio. For example, the coupling constant $g_{a_0NN}$ varies between...
≈ 3.11 and ≈ 10 depending on the $NN$ potential model version \[20\]. Because of that we give below the values for the gradient and pseudoscalar $\eta N$ couplings following from the $a_0\pi N$ triangular $\eta N$ vertex as a function of the $a_0 N$ coupling constant:

$$|G_1(k^2 = m_{\eta}^2)| = 0.06 g_{a_0NN}, \quad |G_2(k^2 = m_{\eta}^2)| = 0.22 g_{a_0NN}.$$  (37)

There are the quantities in Eqs. \(34\) and \(35\) which we shall interpret as the effective pseudovector and pseudoscalar $\eta N$ coupling constants, respectively,

$$f_{\eta NN}^{\text{eff}} = G_1(k^2), \quad g_{\eta NN}^{\text{eff}} = G_2(k^2).$$  (38)

Data analyses on $\eta$ photoproduction off proton near threshold suggest for the pseudoscalar $\eta N$ coupling the small value of $g_{\eta NN}^2/4\pi \approx 0.4$. The value of $g_{a_0NN}$ that fits this number corresponds to the maximal magnitude of $g_{a_0NN}^2/4\pi \approx 6.79$ reported in \[20\]. From Eq. \(31\) one sees that the triangular $a_0\pi N$ correction to the $\eta NN$ vertex represents a mixture \[22\] of pseudovector and pseudoscalar types of $\eta N$ couplings. This mixing is quite important for reproducing the form of the differential cross section for $\eta$ photoproduction off proton at threshold in Fig. 5. Note that such a mixing cannot take place for Goldstone bosons because their point like gradient couplings to quarks are determined in an unique way. On the contrary, in case of extended effective meson-nucleon vertices, such a mixing can take place by means of Eq. \(31\). Data compatibility with the PV-PS mixing created by the $\pi a_0 N$ triangular correction to the $\eta NN$ vertex is a further hint on the non–octet Goldstone boson nature of the $\eta$ meson \[4\].

\[4\]It should be noted that the analytical expressions for the $\eta N$ couplings in Eqs. \(37\) and \(34\) differ from those obtained in \[23\] where the ambiguity in treating the off–shellness of the outgoing proton was not kept minimal. In the present calculation, the on–shell approximation $p \cdot k = -p' \cdot k = -m_{\eta}^2/2$ was made only in evaluating the denominators of the Feynman diagrams, whereas in the nominators $p'$ was consequently replaced by $p' = \hat{p} + \frac{k}{2}$. In contrast to this, in \[23\] the above on–shell
V. EFFECTIVE $F_1(1420)NN$ VER�ICES

The internal structure of the axial vector meson $f_1(1420)$ is still subject to some debates (see Note on $f_1(1420)$ in [19]). Within the constituent quark model this meson is considered as the candidate for the axial meson $(\bar{s}s)$ state and therefore as the parity partner to $\phi$ from the vector meson nonet. The basic difference between the neutral $1^-$ and $1^+$ vector mesons is that while the physical $\omega$ and $\phi$ mesons are almost perfect non-strange and strange quarkonia, respectively, their corresponding parity partners $f_1(1285)$ and $f_1(1420)$ are not. For these axial vector mesons strange and non–strange quarkonia appear mixed up by the angle $\epsilon \approx 15^\circ$:

$$|f_1(1285)\rangle = -\sin \epsilon \langle |\bar{s}s\rangle + \cos \epsilon \frac{1}{\sqrt{2}}(|\bar{u}u + \bar{d}d\rangle, \quad (39)$$

$$|f_1(1420)\rangle = \cos \epsilon \langle |\bar{s}s\rangle + \sin \epsilon \frac{1}{\sqrt{2}}(|\bar{u}u + \bar{d}d\rangle, \quad \epsilon \approx 15^\circ. \quad (40)$$

Therefore within this scheme the violation of the OZI rule for the neutral axial vector mesons appears quite different as compared to the pseudoscalar mesons where the angle corresponding to $\epsilon$ was found to be $\epsilon \approx -45^\circ$.

In other words, while the OZI rule is respected by the internal flavor structure of the vector mesons (there is an almost complete separation between the strange and non–strange quarkonia), it is violated for the axial ones. The latter effect parallels the situation within the $0^-$ octet and is interpreted as the consequence of the $U(1)_A$ anomaly, a subject discussed in [14]. On the other side, the $f_1(1420)$ meson seems alternatively to be equally well interpreted as a $K^*\bar{K}$ molecule [25]. The coupling of this meson to the nucleon is not experimentally well established so far. The only information about it can be obtained from fitting $NN$ approximation was applied to the nominators too and, in addition, terms containing $p'$ have been occasionally interpreted as independent couplings. Through the improper treatment of the off-shellness of the outgoing proton in [23], logarithmically divergent integrals have been artificially invoked in $g_{\eta NN}$.  

17
phase shifts by means of generalized boson exchange potentials of the type considered in Refs. [20,21]. There, one finds that the coupling $g_{f_1NN} \approx 10$ to the nucleon of an effective $f_1$ meson having same mass as the $f_1(1285)$ meson is comparable to the $\omega N$ coupling. Below we demonstrate that couplings of that magnitude can be associated to a large amount with triangular diagrams of the type $KK^*Y$.

To calculate the $KK^*Y$ triangular coupling in (Fig. 4) we use the following effective lagrangians:

\[
\mathcal{L}_{f_1KK^*}(x) = f_{f_1KK^*} \frac{m_{K^*}^2 - m_K^2}{m_K} K^\dagger(x)K^*(x)^\mu f_1(x)_\mu + h.c.,
\]

\[
\mathcal{L}_{KYN}(x) = g_{KNY} i \bar{Y}(x) \gamma_5 N(x) K(x) + h.c.,
\]

\[
\mathcal{L}_{K^*YN}(x) = -g_{K^*YN} \bar{N}(x)(\gamma^\mu K^*(x)_\mu + \frac{\kappa_V}{m_N + m_Y} \sigma^{\mu\nu} \partial_\nu K^*(x)_\mu)Y(x) + h.c.
\]

Here, $g_{KNY}$ stands for the pseudoscalar coupling of the kaon to the nucleon-hyperon ($Y$) system, $g_{K^*NY}$ denotes the vectorial $K^*NY$ coupling, $g_{K^*NY} \frac{\kappa_V}{m_N + m_Y}$ is the tensor coupling of the $K^*$ meson to the baryon, $K^*(x)_\mu$, and $f_1(x)_\mu$ in turn stand for the polarization vectors of the $K^*$ and $f_1$ mesons, $K(x)$ is the kaon field, while $N(x)$ and $Y(x)$ are in turn the nucleon and hyperon fields. A comment on the construction of the effective lagrangian $\mathcal{L}_{f_1KK^*}$ is in place. Ogievetsky and Zupnik (OZ) [26] constructed a chirally invariant lagrangian containing no more than two field derivatives to describe the dynamics of the $a_1(1260)\rho(770)\pi$ system. Exploiting the fact that these mesons have the same external quantum numbers $J^{PC}$ as the respective $f_1(1420)$, $K^*$, and $K$ mesons, the OZ lagrangian might serve as a model for the $f_1K^*K$ system. However, the calculation shows up UV-divergent integrals which are to be maintained by additional cut-offs, i.e. by introducing form factors at both the $KYN$- and $K^*YN$-vertices [27]. This inconvenience makes the lagrangian choice according to [26] less attractive than the present one given above by Eq. (41).

We adopt for the coupling constants the values implied by the Jülich potential [18]

\[
\frac{g_{pKK^+}}{4\pi} = (-0.952)^2, \quad \frac{g_{pKK^*}}{4\pi} = (-1.588)^2, \quad \kappa_V = 4.5.
\]

The value for $f_{f_1KK^*} = 1.97$ has been extracted from the experimental $\Gamma_{K^*+h.c.}$ width of
17 MeV [19]. The amplitude \( T_{f_1N}(K + N \to K^* + N) \) entering the effective \( f_1NN \) vertex can now be expanded into the complete set of the following invariants

\[
T_{f_1N}(K + N \to K^* + N) = \bar{U}_N(p') \left( F_1(k^2) \gamma \cdot \epsilon_{f_1} \gamma_5 + \frac{F_2(k^2)}{m_{f_1}} \gamma \cdot \epsilon_{f_1} \gamma_5 \gamma \cdot k \right) + \frac{F_3(k^2)}{m_Nm_{f_1}} p \cdot \epsilon_{f_1} \gamma_5 \gamma \cdot k + \frac{F_4(k^2)}{m_N} p \cdot \epsilon_{f_1} \gamma_5 \gamma \cdot k \bar{U}_N(p) .
\] (45)

It will become clear in due course that the vector part of the \( K^*NY \) coupling contributes only to the \( F_2(k^2) \) and \( F_4(k^2) \) currents in Eq. (45). All the remaining terms are entirely due to the \( K^*NY \) tensor coupling. For small external momenta, i.e. when \( \not{p}'U_N(p) \approx m_NU_N(p) \), the number of the invariants reduces to three as the first term in Eq. (45) can be approximated by the following linear combinations of the \( F_2(k^2) \) and \( F_3(k^2) \) currents:

\[
m_N\bar{U}_N(p')\gamma_5\gamma_\mu U_N(p) = -p' \cdot k \bar{U}_N(p')\gamma_5\gamma_\mu U_N(p) + p'_\mu \bar{U}_N(p')\gamma_5 \not{k} \bar{U}_N(p) .
\] (46)

One can make use of the approximation of Eq. (46) to estimate \( F_1(k^2) \), the only invariant function which cannot be deduced in an unique way from the triangular \( KK^*Y \) diagrams because of the (logarithmically) divergent parts contained there.

The invariant functions \( F_i(k^2) \) are now evaluated by means of the diagrams of momentum–flow in Fig. 3 in using the techniques of the previous section. The various \( f_1N \) couplings can then be expressed in terms of the integrals

\[
F_i(k^2) = \frac{1}{16\pi^2} f_{f_1KK^*g_{KY}g_{K^*Y}} \int_0^1 dx \int_0^1 dy B_i(k^2) .
\] (47)

The only divergent analytical expression is the one for \( B_1(k^2) \)

\[
B_1(k^2) = 2k \frac{m_{f_1}^2 - m_{K^*}^2}{m_K} m_N(xm_Y - 2m_N) \frac{1}{Z} \\
+ \bar{k} k^2 \frac{m_{f_1}^2 - m_{K^*}^2}{m_K} (xy + (1 - x)(8xy + x - 6)) \frac{1}{Z} , \\
- 2k \frac{m_{f_1}^2 - m_{K^*}^2}{m_K} \ln \frac{Z(m_Y, m_K, \Lambda_{K^*}, x, y, k^2)Z(m_Y, \Lambda_K, m_{K^*}, x, y, k^2)}{Z(m_Y, m_K, m_{K^*}, x, y, k^2)Z(m_Y, \Lambda_K, m_{K^*}, x, y, k^2)} .
\] (48)

All the remaining invariants are convergent and given below as
\[ B_2(k^2) = \frac{m_{f_1}^2 - m_{K^*}^2}{m_K} m_{f_1} [\bar{\kappa} (m_Y (-x^2 y(1-y) + 3xy - 3x + 2) + xym_N) + (1 - x(1+y))] \frac{1}{Z}, \]

\[ B_3(k^2) = -\bar{\kappa} m_{f_1} m_N \frac{m_{f_1}^2 - m_{K^*}^2}{m_K} ((1-y)(x^2(1+3y) - 3x) + 2) \frac{1}{Z}, \]

\[ B_4(k^2) = -2m_N \frac{m_{f_1}^2 - m_{K^*}^2}{m_K} [(x(1-y) - 1)(1 + \bar{\kappa} m_N x(1-y)) + \bar{\kappa} m_Y x(1+y)] \frac{1}{Z}, \]

\[ Z = Z(m_Y, m_K, m_{K^*}, x, y, k^2). \]

(49)

The numerical evaluation of the expressions in the last two equations leads to the following results:

\[ F_1(k^2 = m_{f_1}^2) = -8.56, \quad F_2(k^2 = m_{f_1}^2) = 6.83, \]
\[ F_3(k^2 = m_{f_1}^2) = -3.77, \quad F_4(k^2 = m_{f_1}^2) = -2.03, \]
\[ m_{f_1} = 1385.7 \text{ MeV}. \]

(50)

These results show that all the couplings in Eq. (45) are significant and have to be taken into account in calculating \( \eta \) and \( f_1 \) meson production cross sections.

VI. SUMMARY

As soon as we take three flavor symmetry to emerge from \( U(4)_F \) restricted to \( SU(2)_{ud} \otimes SU(2)_{cs} \otimes U(1) \) in the limit of ‘frozen’ charm degree of freedom, we showed that the hypercharge axial current is anomalous, that the \( \eta \) meson acts as a ‘would be’ strange Goldstone boson, and that the contact \( \eta N \) and \( f_1 N \) couplings proceed via a purely strange isosinglet axial current. For this reason, the couplings considered appeared proportional to \( \Delta s \), the small fraction of nucleon helicity carried by the strange quark sea and strongly suppressed relative quark model predictions. Within this scenario loop vertex corrections acquired importance. We used effective lagrangians to construct effective \( \eta N \) and \( f_1 N \) coupling strengths in terms of triangular \( a_0 \pi N \), and \( KK^*Y \) diagrams, respectively. We found all kinds of effective couplings (up to one) to be determined by divergenceless expressions. This is the reason for which we view the model developed here as useful for describing \( \eta N \) and \( f_1(1420)N \) couplings beyond the limits of applicability of ChPT. The strengths of all
the invariant amplitudes contributing to the spatially extended triangular meson–nucleon vertices were found to be of significant size and the produced mixing between the different types of meson–nucleon couplings was shown to be important for data interpretation. Especially the form of the differential cross section for η photoproduction off proton at threshold was successfully reproduced in terms of a mixing between pseudoscalar and pseudovector ηN couplings brought about by the invariant decomposition of the $a_0 + N \rightarrow \pi + N$ scattering amplitude entering the triangular $a_0\pi N$ vertex.

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Figures

Fig. 1: The η pole term.

Fig. 2: The triangular ηNN vertex.

Fig. 3: The flow of momentum.

Fig. 4: The triangular $f_1NN$ vertex.

Fig. 5: Differential cross section for η photoproduction off proton at lab energy of 724 MeV as calculated within the model of Tiator, Bennhold and Kamalov [16]. The dotted and dash–dotted lines correspond to $g_{\eta NN}^2/4\pi$ taking the values of 0.4 and 1.1, respectively. The full line corresponds to the $a_0\pi N$ triangular coupling. The data are taken from Krusche et al. [16].
REFERENCES

[1] S. Coleman, *Aspects of Symmetry*, Selected Erice Lectures (Cambridge University Press, 1985).

[2] J. Gasser and H. Leutwyler, Nucl. Phys. B291 (1985) 465.

[3] V. Bernard, N. Kaiser and U.-G. Meißner, Int. J. Mod. Phys. E4 (1995) 193.

[4] M. Kirchbach and H.-J. Weber, Comm. Nucl. Part. Phys. 22 (1998) 171.

[5] M. Kirchbach. Phys. Rev. D (1998) (in press); nucl-th/9801030.

[6] M. J. Savage and J. Walden, Phys. Rev. D55 (1997) 5376.

[7] J. Speth and A. W. Thomas, Adv. Nucl. Phys. (in press);
   H. Holtman, A. Szczurek, and J. Speth, Nucl. Phys. A596 (1996) 631.

[8] V. De Alfaro, S. Fubini, G. Furlan, and C. Rossetti, *Currents in Hadron Physics* (Amsterdam: North Holland, 1973).
   R. K. Bhaduri, *Models of the Nucleon* (California, Addison–Wesley, 1988).

[9] R. L. Jaffe, Phys. Lett. B229 (1989) 275.

[10] G. Veneziano, Mod. Phys. Lett. A4 (1989) 1605.

[11] S. A. Coon and M. D. Scadron, Phys. Rev. C42 (1990) 2256.

[12] J. Ambjørn, J. Greensite, and C. Peterson, Nucl. Phys. B221 (1983) 381.

[13] G. ’t Hooft, Phys. Rev. D14 (1976) 3432.

[14] S. Forte and E. Shuryak, Nucl. Phys. B367 (1991) 153.

[15] J. Ellis and M. Karliner, Phys. Lett. B341 (1995) 395.

[16] B. Krusche et al., Phys. Rev. Lett. 74 (1995) 3736;
   L. Tiator, C. Bennhold and S. Kamalov, Nucl. Phys. A580 (1994) 455.
[17] W. Grein and P. Kroll, Nucl. Phys. A338 (1980) 332.

[18] A. Reuber, K. Holinde, and J. Speth, Nucl. Phys. A570 (1991) 541.

[19] Review of Particle Properties, Phys. Rev. D54 (1996) 1.

[20] R. Machleidt, Adv. Nucl. Phys. 19 (1989) 189.

[21] R. Machleidt, K. Holinde, and C. Elster, Phys. Rep. 149 (1987) 149.

[22] F. Gross, J. W. Van Orden, and K. Holinde, Phys. Rev. C41 (1990) R1909.

[23] M. Kirchbach and L. Tiator, Nucl. Phys. A604 (1996) 385.

[24] T. Bolton et al., Phys. Lett. B279 (1992) 495.

[25] A. Reuber, K. Holinde, H. C. Kim, and J. Speth, Nucl. Phys. A608 (1996) 243.

[26] V. I. Ogievetsky and B. M. Zupnik, Nucl. Phys. B24 (1970) 612.

[27] S. Neumeier, Diploma Thesis, TH Darmstadt, Germany, 1996, unpublished.
Fig. 1
Fig. 3
Fig. 4

\[ \Lambda, \Sigma \rightarrow K^* \rightarrow f_1(1420) \]
Fig. 5

724 MeV $\gamma + p \rightarrow \eta + p$

$\frac{d\sigma}{d\Omega}$ (µb/sr)

$\theta_{c.m.}$

0 30 60 90 120 150 180