Meson-Baryon Couplings Revisited

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Abstract

The theoretical evaluation of the coupling constants $g^{\pi NN}$, $g_{KNN}$, and $g_{KNS}$ is undertaken using QCD sum rules. These quantities were previously calculated with exponential (Borel) kernels used to suppress the unknown contributions of the hadronic continua. This method however introduces arbitrariness and instability in the calculation. In order to avoid these I redo the calculation using polynomial kernels tailored to vanish at the baryonic resonance masses. The results are $g^{\pi NN} = 11.5 \pm 1.0$, $g_{KNS} = 5.45 \pm 0.4$, $g_{KNS} = -(12.7 - 15.0)$ which are close to experiment and to the predictions of SU(3) and which, with the corresponding Goldberger-Treiman Discrepancy satisfy quite well the Dashen-Weinstein relation.

1 Introduction

The $\pi N$ coupling constant $g^{\pi NN}$ is one of the most important parameters in hadron physics. This quantity has been extensively studied since the works of Reinders, Rubinstein and Yazaki [1] using the method of QCD sum rules [2] starting with either the correlator of three interpolating fields or with the pion to vacuum matrix element of two nucleon interpolating currents. These correlators introduce several different Dirac structures and these were studied in detail two decades ago [3], [4], [5] and more recently in [6]. All these calculations however involve dispersion integrals which include the unknown contribution of the continuum with quantum numbers of the nucleon. In order to minimize this unknown contribution a damping kernel $e^{-M^2}$ is introduced. This method however presents problems of arbitrariness and stability. I present here a different choice for the damping kernel: a polynomial $P(t)$ which practically vanishes in the nucleon resonance region. The sum rule method consists of expressing the quantity of interest, which enters in the residue at the pole of the correlation function $\Pi(t)$ in terms of an integration of $\Pi(t)$ over the contour $c$ shown in Fig.1 in the complex t-plane. This in turn is the sum of an
integral over the cut of the discontinuity of $\Pi(t)$ and an integral over the circle of large radius $R$ over which $\Pi(t)$ can be replaced by its QCD expression $\Pi^{QCD}(t)$ the latter can be brought back to an integral over the cut of the discontinuity of $\Pi^{QCD}(t)$. The main unknown in the calculation is the integral over the cut of the hadronic amplitude.

In order to minimize this contribution an integration kernel $P(t)$ is introduced so that

$$\text{Residue } x P(\text{pole}) = \frac{1}{\pi} \int_{\text{th}}^{R} dt P(t) \text{Im} \Pi(t) + \frac{1}{\pi} \int_{0}^{R} dt P(t) \text{Im} \Pi^{QCD}(t)$$

(1)

The usual choice for the kernel is $P(t) = e^{-tM^2}$ where $M^2$ (the Borel mass) is a damping parameter which cannot be too large because the damping worsens nor can it be too small because the contribution of higher orders in the Operator Product Expansion (OPE) of $\Pi^{QCD}(t)$ increase beyond control. An intermediate value of $M^2$ has to be chosen by stability considerations which is not always possible. The radius $R$ of the circle is another adjustable parameter.

In order to avoid these problems I shall choose for the kernel a simple polynomial $P(t) = 1 - a_1 t - a_2 t^2$ tailored to practically vanish in the nucleon resonance region in order to render the first integral on the r.h.s. of eq. (1) negligible. The value of $R$ is chosen to lie in the stability region of the second integral on the r.h.s. of eq. (1). This leads to an unambiguous result.

The method is applied to the calculation of the coupling constants $g_{\pi NN}, g_{KN\Lambda}$ and $g_{KN\Sigma}$.

### 2 The Pion-Nucleon Coupling Constant $g_{\pi NN}$

Consider the correlation function

$$\Pi(p, q) = i \int dx e^{iqx} \langle 0 | T j_p(x) j_n(0) | \pi^+(p) \rangle$$

(2)
where
\[ j_p = \epsilon_{abc} (u_a^T C \gamma_\mu u_b) \gamma_5 \gamma_\mu d_c \]
(3)
is the proton interpolating field of Joffe [7]. The interpolating field of the neutron
is obtained by the interchange of \( u \) and \( d \).

I shall work with the \( \gamma_5 \sigma_{\mu\nu} q_\mu q_\nu \) structure of the correlator.

In the soft pion limit the OPE reads [5]:
\[ \Pi_{QCD}(t) = \sqrt{2} \langle qq \rangle \left[ \frac{\ln(-t)}{12\pi^2 f_\pi} + \frac{4 f_\pi}{3 t} - C \right] \]
(4)
with \( t = q^2 \),
\[ C \simeq \left( \frac{1}{216 f_\pi} \langle a_s GG \rangle - \frac{m_N^2 f_\pi}{6} \right) \simeq -0.0117 \text{GeV}^3 \]
(5)
at the standard values \( \langle a_s GG \rangle \approx 0.012 \text{GeV}^4, m_N^2 \approx 0.8 \text{GeV}^2 \).

On the hadronic side \( \Pi(t) \) has a double nucleon pole, single nucleon poles
and a continuum.
\[ \Pi(t) = \sqrt{2} g_{\pi NN} \lambda_N^2 \frac{(t - m_N^2)}{(t - m_N^2)^2} + \text{single poles} + \text{continuum} \]
(6)
\( \lambda_N \) denotes the coupling of \( j_N \) to the nucleon
\[ \langle 0 | j_N | N \rangle = \lambda_N U_N \]
(7)

Consider the integral \( \frac{1}{2\pi i} \int_c dt P(t) \Pi(t) \) in the complex t-plane where \( c \) is the contour shown in Fig.1.

As argued in the introduction I shall choose \( P(t) = 1 - a_1 t - a_2 t^2 \) with the coefficients \( a_{1,2} \) tailored in order to practically annihilate the contribution of the continuum. We thus have
\[ \sqrt{2} g_{\pi NN} \lambda_N^2 P'(m_N^2) + \Delta = \frac{1}{\pi} \int_{i\theta}^R dt P(t) \text{Im} \Pi(t) \Pi + \frac{1}{2\pi i} \oint dt P(t) \Pi_{QCD}(t) \]
(8)
\( \Delta \) is the contribution of the single poles. The choice of \( P(t) \) is based on
the fact that the contribution of the first integral on the r.h.s. of eq. (7) arises
mostly from the interval [8] \( I = 2.0 \text{GeV}^2 \leq t \leq 3.0 \text{GeV}^2 \) where the resonances
\( N^+ (1440), N^- (1535), N^- (1650), N^+ (1710) \) lie. The parameters \( a_1 \) and \( a_2 \) are
chosen so as to minimize the integral \( \int_{2.0 \text{GeV}^2}^{3.0 \text{GeV}^2} dt \ |P(t)|^2 \), numerically \( a_1 = .807 \text{GeV}^{-2}, a_2 = -.160 \text{GeV}^{-4} \).

This allows the neglect of the first integral on the r.h.s. of eq. (8) so that, using expression (4) for \( \Pi_{QCD}(t) \)

\[ \sqrt{2} g_{\pi NN} \lambda_N^2 P'(m_N^2) + \Delta \simeq \sqrt{2} \langle qq \rangle \left[ \frac{1}{12\pi^2 f_\pi} \text{I}_0(R) + \frac{4 f_\pi}{3} - a_1 C \right] \]
(9)
with the definition
\[ I_n(R) = \int_0^R dt t^n P(t) \]  
(10)

The fact that the contribution of term \( a_1 C \) is small compared to that of the preceding ones (\( \sim 5\% \)) justifies the neglect of the next unknown term proportional to \( a_2 \). This is to be contrasted to the case of exponential damping where an infinite number of unknown higher order terms enter. Moreover it can be seen from Fig. 2 that the damping of the continuum is much better in our case.

\[ P(t) \text{ is our polynomial } 1 - a_1 t - a_2 t^2 \]

\[ M^2 \sim 1.0 GeV^2 \text{ as chosen in the literature} \]

\( \lambda_N \) was obtained using the same approach in [8] starting with the nucleon correlator
\[ \Pi_N(q) = i \int dx e^{ix\mu} \langle 0 | T \eta(x) \eta(0) | 0 \rangle = q \Pi_1(q^2) + \Pi_2(q^2) \]  
(11)

The sum rule using \( \Pi_2(q^2) \) yields
\[ (2\pi)^4 m_N^2 \lambda_N^2 P(m_N^2) = -B_3 I_1(R) - B_7 \]  
(12)

with the OPE coefficients
\[ B_3 = 4\pi^2 (1 + \frac{3}{2} a_s) \langle \bar{q}q \rangle, \quad B_7 = -\frac{4}{3} \pi^2 \langle \bar{q}q \rangle \langle a_s GG \rangle \]  
(13)

Let us now estimate \( \Delta \). The contribution of the \( N'(1440) \) adds to eq. (6)
\[ \delta_{N'} = \sqrt{2} g_{\pi NN'} \lambda_N \lambda_{N'} \frac{1}{(t - m_N^2)(t - m_{N'}^2)} \]  
(14)

which corresponds to
\[ \Delta = \sqrt{2} g_{\pi NN} \lambda_N \lambda_{N'} \frac{P(m_N^2) - P(m_{N'}^2)}{m_N^2 - m_{N'}^2} \]  
(15)
While \( P(m^2_{N'}) \) is small by construction \( P(m^2_N) \) is not. The contribution of the simple pole is not damped by the polynomial. 

\( \lambda_N \) is estimated from a variant of eq. (12) (with no damping)

\[
(2\pi)^4 \lambda_N^2 m_{N'} \simeq -B_3 \int_{(m^2_N + m^2_{N'})/2}^{(m^2_N + m^2_{N'})/2} t \, dt 
\]

which yields \( \frac{\lambda_N}{\lambda_N'} \simeq .35 \)

The coupling constant \( g_{\pi NN'} \) was studied in [9] who give \( g_{\pi NN'} = .47 \pm .04 \). This yields

\[
\frac{\Delta_{NN'}}{\text{dominant term}} = g_{\pi NN'} \frac{\lambda_{NN'}}{\lambda_N} \frac{P(m^2_N)}{P(m^2_{N'})} \simeq .015 \quad \text{which shows that } \Delta \text{ amounts to only a few percent of the dominant term. Finally}
\]

\[
g_{\pi NN} - P'(m^2_N) \frac{P(m^2_{N'})}{P(m^2_N)} = \left[ \frac{1}{4\pi^2} I_0(R) + \frac{1}{2} f_\pi^2 - a_1 f_\pi C \right] \frac{m_N}{f_\pi} \frac{m_N}{f_\pi} \]

As discussed before the radius \( R \) of the integration circle is not arbitrary, it cannot be too small because this would invalidate the OPE nor can it be too large because the polynomial would start enhancing the unknown continuum instead of damping it, \( R \) should be chosen in the stability region of the integrals \( I_{0,1}(R) \) shown in Fig. 3

The interval of stability is very wide and the results are

\( I_0(R) = .814 GeV^2 \) similarly \( I_1(R) = .474 GeV^4 \). With the standard value \( \langle a_sGG \rangle = .012 GeV^4 \) eq. (17) finally gives

\[
g_{\pi NN} = 11.5 \pm 1.0
\]

The error, estimated about 8% stems mostly from neglected radiative corrections as well as from \( \Delta \).

In order to assess the stability of the method other choices of the damping polynomial can be made, e.g. \( P(t) = (1 - \frac{t}{2.5 GeV^2}) \) or impose \( \int_{2.0 GeV^2}^{3.0 GeV^2} dt \, P(t) = 0 \). In both cases the results are very close to eq. (18) \( (g_{\pi NN} = 11.2 \) and \( g_{\pi NN} = 11.0) \)
3 Kaon-Baryon Couplings

Several attempts were made to evaluate the couplings $g_{KN\Lambda}$ and $g_{KN\Sigma}$ using QCD sum rules \[10, 11, 12, 13, 14, 15\] which all turned out much smaller than their SU(3) values

\[ g_{KN\Lambda} = \frac{-1}{\sqrt{3}}(3 - 2\alpha)g_{\pi NN} \simeq -13.05 \]  
\[ g_{KN\Sigma} = (2\alpha - 1)g_{\pi NN} \simeq 3.50 \]  

where $\alpha = D/(D + F) \simeq 0.635$ is the fraction of D-type coupling.

Experimentally, the determination of the coupling constants still involves uncertainties but the analysis of \[16\] yields

\[ g_{KN\Lambda} = -13.5 \quad \text{and} \quad g_{KN\Sigma} = 4.5 \]  

Evaluation of the coupling constants using heavy baryon chiral perturbation theory at the one loop level has also been attempted \[17\] resulting in values much smaller than one would expect from eq. (19).

An attempt to calculate the coupling constants starting from the Golberger-Treiman relation (GTR) in $SU(3) \times SU(3)$ was undertaken in \[18\]. In this paper the extrapolation to the nucleon mass shell was done using exponential kernels to damp the unknown parts of dispersion integrals. I shall here redo the calculation using polynomial kernels which avoid the arbitrariness and stability problems.

The GTRs in the strange channel are

\[ \sqrt{2}f_K g_{KN\Lambda} = (m_N + m_Y)g_A^Y(0) \]  

where $Y = \Lambda, \Sigma$, $g_A^Y(0)$ is the axial-vector coupling constant and $f_K$ the K-meson decay constant.

The coupling constants could of course be obtained from eq. (21) because the $g_A^Y$ are measured. eq. (21) involves however an extrapolation in momentum transfer squared from 0 to $m_K^2$ which is not a small quantity on the hadronic scale. In other words explicit chiral symmetry breaking leads to corrections to the GTR, the GT discrepancies (GTD) which are not small in the strangeness changing case.

The GTD are defined by

\[ \Delta^Y = 1 - (m_N + m_Y)g_A^Y(0)/(\sqrt{2}f_K g_{KN\Lambda}) \]  

The evaluation of the K-nucleon coupling constants will yield the GTD.

Start from the matrix element

\[ \langle P \left| \partial_\mu A^{K^+}_\mu \right| Y \rangle = \Pi(q^2)\not{q} \gamma_5 Y \]  

$q$ denoting the momentum transfer between the baryons. We have
\[ \Pi(0) = (m_N + m_Y)g_A(0) \] (24)

The analytic properties of \( \Pi(t = q^2) \) in the complex \( t \)-plane are known. It has a pole at \( t = m_K^2 \) and a cut along the positive \( t \)-axis starting at \( t_{th} = (m_K + 2m_\pi)^2 \)

\[ \Pi(t) = -\sqrt{2}f_Km_K^2 g_{KNY} \frac{g_{KN\pi}}{(t - m_K^2)} + \cdots \] (25)

Consider now the usual contour \( c \) in the complex \( t \)-plane and the integral

\[ \frac{1}{2\pi i} \int_c \frac{dt}{t} \Pi(t) = \frac{1}{\pi} \int_{th}^R \frac{dt}{t} \text{Im}\Pi(t) + \frac{1}{2\pi i} \oint \frac{dt}{t} \Pi^{QCD}(t) = \Pi(0) - \sqrt{2}f_Kg_{KNY} \] (26)

The first integral on the r.h.s. running along the cut represents the contribution of the \( 0^- \) strange continuum and provides the main part of the GTD. In the second integral along the circle \( \Pi^{QCD}(t) \) is a good approximation for \( \Pi(t) \). In order to overcome the lack of knowledge of \( \text{Im}\Pi(t) \) on the cut consider the modified integral

\[ \frac{1}{2\pi i} \int_c dt \left( \frac{1}{t} - a_0 - a_1t \right) \Pi(t) = \frac{1}{\pi} \int_{th}^R dt \left( \frac{1}{t} - a_0 - a_1t \right) \text{Im}\Pi(t) + \frac{1}{2\pi i} \oint dt \left( \frac{1}{t} - a_0 - a_1t \right) \Pi^{QCD}(t) \] (27)

The main contribution to the integral over the cut arises from the interval \( I : 1.5 GeV^2 \lesssim t \lesssim 3.5 GeV^2 \) which includes the resonances \( K(1400) \) and \( K(1830) \). The constants \( a_1 \) and \( a_2 \) are now chosen so as to annihilate the kernel \( (1 - a_1t - a_2t^2) \) at \( t = 1.46^2 GeV^2 \) and \( t = 1.83^2 GeV^2 \) i.e

\[ a_0 = 0.77 GeV^{-2} \quad \text{and} \quad a_1 = -0.14 GeV^{-4} \] (28)

With this choice the integrand is reduced to only a few percent of its initial value over the interval \( I \) and the integral over the cut becomes negligible. so

\[ \Pi(0) - \sqrt{2}f_Kg_{KNY}(1 - a_1m_K^2 - a_2m_K^4) \simeq \frac{1}{2\pi i} \oint dt \left( \frac{1}{t} - a_0 - a_1t \right) \Pi^{QCD}(t) \] (29)

It appears from the equation above that chiral symmetry breaking manifests itself in the presence of the r.h.s. as well as in the deviation of the factor \( (1 - a_1m_K^2 - a_2m_K^4) \) from unity.

In [18] \( \Pi^{QCD}(t) \) was obtained from Borel-type sum rules. I shall here redo the calculation using a polynomial kernel as advocated. Consider the three-point function
\[
\Gamma(s = p^2, t = q^2) = \int \int \, dx \, dy \, e^{-ipx} \, e^{iqy} \langle 0 | T\psi^N(x) \partial_\mu A^K_\mu(y) \psi^Y(0) | 0 \rangle \tag{30}
\]

where \(\psi^N,Y\) are baryonic currents and let

\[
\Gamma(s, t) = F(s, t) \sigma_{\mu \nu} \gamma_5 q_\mu p'_\nu + \text{other tensor structure} \tag{31}
\]

Furthermore the residue at the double baryonic pole of the \(\Gamma\) is related to \(\Pi(t)\), i.e.

\[
\Gamma(s, t) = \left( \frac{\lambda_N \lambda_Y \Pi(t)}{(s - m_N^2)(s - m_Y^2)} + \cdots \right) \sigma_{\mu \nu} \gamma_5 p'_\mu q'\nu + \cdots \tag{32}
\]

\(\Pi^{QCD}(t)\) is obtained from the above by using \(F^{QCD}(s, t)\), extrapolating to the baryons mass shell and identifying terms in \(\Pi(t)\).

\(\lambda_{N,Y}\) denotes the coupling of the baryonic currents to the corresponding baryons.

It will be assumed that the contribution of the single poles is small as it was in the \(\pi - N\) case and it will be neglected.

\(\Pi^{QCD}(t)\) is obtained by using \(F^{QCD}(t)\) and extrapolating to the baryon mass shell.

\(F^{QCD}(s, t)\) is given in \(\Pi^{QCD}(s, t)\)

\[
F^{QCD}_{\Lambda}(s) = \frac{C_{\Lambda}}{s t}, \quad C_{\Lambda} = \frac{4}{3} \sqrt{\frac{2}{3}} m_s (\langle qq^2 \rangle + \langle q\bar{q} \bar{q}s \rangle) \tag{33}
\]

and

\[
F^{QCD}_{\Sigma}(s, t) \cong 0 \tag{34}
\]

Consider now the integral

\[
\frac{1}{2\pi i} \int \, ds P_Y(s) F(s, t) \tag{35}
\]

in the complex \(s\)-plane where as before \(P_Y(s)\) is the damping polynomial intended to eliminate the contribution of the continuum.

This gives

\[
\lambda_N \lambda_Y \Pi^{QCD}(t) \left( \frac{P_Y(m_N^2) - P_Y(m_Y^2)}{m_N^2 - m_Y^2} \right) = \frac{1}{2\pi i} \int \, ds P_Y(s) F^{QCD}(s) \tag{36}
\]

and take \(P_Y(s) = (1 - \frac{s}{m_N^2})(1 - \frac{s}{m_Y^2})\) which vanishes at the first nucleon and hyperon excited states \((m_N^2 = 1.98 \text{GeV}^2, \ m_Y^2 = 2.40 \text{GeV}^2, \ m_{\Lambda'}^2 = 2.51 \text{GeV}^2)\).

Eqs. \(33\) and \(34\) inserted in eq. \(35\) yield the QCD expressions

\[
\frac{\lambda_N \lambda_{\Lambda}}{(m_{\Lambda'}^2 - m_N^2)} (P_{\Lambda}(m_{\Lambda'}^2) - P_{\Lambda}(m_N^2))) \Pi^{QCD}(t) = \frac{C_{\Lambda}}{t} \tag{37}
\]

and
\[ \lambda_N \lambda_\Sigma \left( \frac{m_\Sigma^2 - m_N^2}{m_\Lambda^2 - m_\Sigma^2} \right) \left( \frac{P_\Sigma(m_\Sigma^2) - P_\Sigma(m_N^2)}{P_\Sigma(m_\Lambda^2) - P_\Sigma(m_\Sigma^2)} \right) \Pi_{\Sigma}^{QCD}(t) \geq 0 \]  

(38)

When this is inserted in eq. (29) it gives

\[ \Pi_\Lambda(0) = \sqrt{2} f_K g_{K \Lambda \Sigma}(1 - a_1 m_K^2 - a_2 m_K^4) - \frac{a_1 C_\Lambda(m_\Lambda^2 - m_N^2)}{\lambda_N \lambda_\Sigma \left( \frac{m_\Lambda^2 - m_N^2}{m_\Lambda^2 - m_\Sigma^2} \right) \left( \frac{P_\Lambda(m_\Lambda^2) - P_\Lambda(m_\Sigma^2)}{P_\Lambda(m_\Lambda^2) - P_\Lambda(m_\Sigma^2)} \right)} \]  

(39)

\[ \Pi_\Sigma(0) = \sqrt{2} f_K g_{K \Lambda \Sigma}(1 - a_1 m_K^2 - a_2 m_K^4) \]  

(40)

\[ \lambda_N \] is given by eq. (12)

\[ (2\pi)^4 \lambda_N^2 P(m_N^2) = 0.444 \text{GeV}^6 \]  

(41)

a similar treatment of the \( \Lambda \) 2-point function \[12\] with our polynomial replacing the damping exponential gives

\[ (2\pi)^4 \lambda_\Lambda^2 P_\Lambda(m_\Lambda^2) = 0.481 \text{GeV}^6 \]  

(42)

Finally with the experimental values \( g_\Lambda^A = -0.72 \), \( g_\Sigma^A = 0.34 \) we obtain

\[ g_{K \Sigma \Sigma} = 5.45, \quad g_{K N \Lambda} = -(11.50 \pm 1.17 \kappa) \]  

(43)

\( \kappa \) denotes the deviation of the 4-quark condensate from factorization, it arises from the condensate \( \langle \bar{q} q \bar{q} q \rangle \simeq \kappa \langle \bar{q} q \rangle \langle \bar{q} q \rangle \) which appears in the expression for the constant \( C_\Lambda \).

There is no consensus on the value of \( \kappa \). I shall take \( 1 \leq \kappa \leq 3 \).

4 Results and Conclusions

The \( \pi - N \) and K-N coupling constants were calculated using QCD sum-rules with polynomial kernels tailored to vanish in the baryonic resonance region. This solves the problems of arbitrariness and stability inherent to the usual Borel-type QCD sum-rules. The results are

\[ g_{\pi NN} = 12.4 \pm .6, \quad g_{K \Sigma \Sigma} = 5.45 \pm .4, \quad g_{K N \Lambda} = -(12.7 - 15.0) \]  

(44)

which are quite close to experiment and to the SU(3) values

The corresponding GTD are large as expected

\[ \Delta^\Sigma_{GT} \simeq .18, \quad \Delta^\Lambda_{GT} \simeq .29 \]  

(45)

It is finally interesting to see how well our results fit the Dashen-Weinstein \[19\] relation between the GTD

\[ g_{\pi NN} \Delta^\Sigma_{GT} = \frac{m_\pi^2}{2m_K^2} (g_{K \Sigma \Sigma} \Delta^\Sigma_{GT} - \sqrt{3} g_{K N \Lambda} \Delta^\Lambda_{GT}) \]  

(46)

With \( \Delta^\Sigma_{GT} \simeq .02 \) it is seen that relation eq. (46) is quite well satisfied.
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