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Strange vector currents and the OZI–rule

Ulf-G. Meißner[‡,1)], V. Mull[‡ #1] , J. Speth[‡,2)], J. W. Van Orden[†,3)]

[‡]Forschungszentrum Jülich, IKP (Theorie), D–52425 Jülich, Germany
[†]Department of Physics, Old Dominion University, Norfolk, VA 23529, USA
and
Jefferson Lab., 12000 Jefferson Ave., Newport News, VA23106, USA

email:
[1)] *Ulf-G.Meissner@kfa-juelich.de,* [2)] *J.Speth@kfa-juelich.de,* [3)] *vanorden@jlab.org*

ABSTRACT

We investigate the role of correlated $\pi\rho$ exchange in the extraction of matrix elements of the strange vector current in the proton. We show that a realistic isoscalar spectral function including this effect leads to sizeably reduced strange vector form factors based on the dispersion–theoretical analysis of the nucleons' electromagnetic form factors.

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[#1] Present address: IBM Informationssysteme GmbH, Köln, Germany.

One of the outstanding problems in the understanding of the nucleon structure concerns the strength of various strange operators in the proton. A dedicated program at Jefferson Laboratory supplemented by experiments at BATES (MIT) and MAMI (Mainz) is aimed at measuring the form factors related to the strange vector current $\bar{s}\gamma_\mu s$ in the nucleon. It was already pointed out a long time ago by Genz and Höhler [1] that the dispersion–theoretical analysis of the nucleons' electromagnetic form factors allows one to get bounds on the violation of the OZI rule, which leads one to expect that strange matrix elements should be small [2]. This rule has, however, never firmly been routed in QCD but can be understood qualitatively in large N_C (with N_C the number of colors) [3]. Jaffe [4] showed that under certain assumptions the information encoded in the isoscalar nucleon form factors can be used to extract strange matrix elements. Of particular importance for this type of analysis is the identification of the two lowest poles in the isoscalar spectral function with the $\omega(782)$ and the $\phi(1020)$ mesons. The corresponding strange form factors turn out to be rather large in magnitude, related to the strong coupling of the ϕ to the nucleon found in the dispersion–theoretical analysis [5]. This analysis was later updated and extended in [6] based on the novel form factor fits presented in [7]. Loosely spoken, such an analysis is based on a "maximal" violation of the OZI rule because the spectral function in the mass region of about 1 GeV is assumed to be given entirely by the ϕ pole. On the other hand, the coupling of various mesons (like the ω and the π) to the nucleons has been investigated in great detail in the framework of the Bonn–Jülich meson exchange potential for the nucleon–nucleon interaction by Holinde and coworkers [8] . In particular, the correlated $\pi\rho$ [8][9] and $\pi\pi$ exchange [10] has recently been included consistently. This leads to a more realistic microscopic picture of the isoscalar spectral function in the mass region of the ϕ [11]. Our aim is to combine this novel results from the NN interaction with a fit to the nucleon electromagnetic form factors and to elucidate the strength of the ϕNN couplings, i.e. the violation of the OZI rule, and the consequences for the extraction of the strange form factors.

To be specific, consider first the nucleon–nucleon interaction. Although QCD is believed to be the theory underlying the strong interactions, in the non–perturbative regime of low– and medium–energy physics, mesons and baryons have retained their importance as effective, collective degrees of freedom for a wide range of nuclear phenomena. This is most apparent for the NN system. Here, the interaction between the two nucleons is generated by meson exchange [12]. Resulting potentials, e.g. the Paris [13], Nijmegen [14] or Bonn [15] potentials, are able to describe the NN data below pion threshold in a truly quantitative manner. The full Bonn potential contains apart from single–meson exchanges higher–order diagrams involving also the Δ–isobar. The strength of the various baryon–meson vertices is parametrized by coupling constants. In addition, form factors with cut-off masses Λ_α are included as additional parameters, they take into account the corresponding vertex extensions. However, there are some longstanding conceptual problems hidden in the choice of parameters which have been resolved in the past years by Holinde and coworkers. First, the fictitious scalar–isoscalar meson $\sigma_{\rm OBE}$, which is needed to provide the intermediate–range attraction, has been replaced by correlated 2π–exchange [10]. A second longstanding discrepancy existed for the cut-off $\Lambda_{\pi NN}$, which is rather large in the present–day potential models ($\sim 1.3\,$GeV) compared to the information from other sources, like e.g. πN scattering. In [8] it has been shown that the interaction between a π and a ρ meson (correlated $\pi\rho$ exchange) has a strong influence on the NN–potential in the pion channel. It provides a sizeable contribution with a peak around 1.1 GeV. Due to this additional π–like contribution one is able to reduce the cut-off $\Lambda_{\pi NN}$, which is now in much better agreement with information from other sources. The third well–known discrepancy is the ωNN coupling constant, which in most of the NN potentials is three times bigger than predicted by SU(3) symmetry [16]. It has also been shown in [9] that in the ω–channel the correlated $\pi\rho$ exchange gives a sizeable contribution which allows the choice of a value for $g_{\omega NN}$ which is in reasonable agreement with the SU(3) prediction. The process missing in the original Bonn potential is depicted in fig. 1. It has been analyzed in detail in [9]. The result is given in terms of a dispersion integral, which for simplicity can be represented by an effective one–boson exchange, denoted as ω', in the ω–channel,

$$\frac{1}{\pi}\int_{(M_\pi+M_\rho)^2}^{\infty}\frac{\rho_P^\omega(s,t')}{t'-t}\,dt' = -\frac{g^2_{\omega' NN}}{t-M^2_{\omega'}} \quad . \tag{1}$$

The spectral function ρ_P^ω is again peaked around 1.1 GeV. The pole fit, eq.(1), gives the following ω' parameters: $M_{\omega'} = 1.12\,$GeV, $g^2_{\omega' NN}/4\pi = 8.5$ for the vector and $f^2_{\omega' NN}/4\pi = 1.5$ for the tensor coupling. Moreover, it turns out that $g_{\omega' NN} < 0$ and $\kappa_{\omega'} = f_{\omega' NN}/g_{\omega' NN} > 0$. In what follows, we use this effective pole instead of the full spectral function. This approach has further been extended to include kaon loops and hyperon excitations, with the parameters fixed from a study of the reactions $p\bar{p} \to \Lambda\bar{\Lambda}$ and $p\bar{p} \to \Sigma\bar{\Sigma}$ [17]. There are sizeable cancellations between the various contributions from graphs (see fig. 2) with intermediate K's, K^*'s and diagrams with the direct hyperon interactions [11] leading to a very small ϕ coupling,

$$\frac{g^2_{\phi NN}}{4\pi} \simeq 0.005\,, \quad \kappa_\phi \simeq \pm 0.2\ . \tag{2}$$

The various contributions are tabulated in table 1. The sign of the tensor coupling is very sensitive to the details of the calculation. The smallness of these couplings amounts to a "resurrection" of the OZI rule. Note that such strong cancellations have also been observed in the quark model study of [18].

The structure of the nucleon as probed with virtual photons is parametrized in terms of the so–called Dirac (F_1) and Pauli (F_2) form factors. These form factors have been measured over a wide range of space–like momentum transfer squared, $t = 0\ldots-35\,\text{GeV}^2$ but also in the time–like region either in $\bar{p}p$ annihilation or in $e^+e^- \to \bar{p}p, \bar{n}n$ collisions. The tool to analyze these data in a largely model–independent fashion is dispersion theory [5].

We therefore briefly review the dispersion–theoretical formalism developed in [7] and discuss the pertinent modifications due to the constraints from the NN interaction described before. Assuming the validity of unsubtracted dispersion relations for the four form factors $F_{1,2}^{(I=0,1)}(t)$ [19], one separates the spectral functions of the pertinent form factors into a hadronic (meson pole) and a quark (pQCD) component as follows,

$$F_i^{(I)}(t) = \tilde{F}_i^{(I)}(t)L(t) = \left[\tilde{F}_i^{\rho}(t)\,\delta_{I1} + \sum_I \frac{a_i^{(I)}\,L^{-1}(M_{(I)}^2)}{M_{(I)}^2 - t}\right]\left[\ln\left(\frac{\Lambda^2 - t}{Q_0^2}\right)\right]^{-\gamma} \tag{3}$$

where $F_i^{\rho}(t) = \tilde{F}_i^{\rho}(t)\,L(t)^{-1}$ parametrizes the isovector ($I = 1$) two–pion contribution (including the one from the ρ) in terms of the pion form factor and the P–wave $\pi\pi\bar{N}N$partial wave amplitudes in a parameter–free manner. In addition, we have three isovector poles, the masses of the first two can be identified with physical ones, i.e. $M_{\rho'} = 1.45\,\text{GeV}$ and $M_{\rho''} = 1.65\,\text{GeV}$. In the isoscalar channel ($I = 0$), we have the poles representing the ω, the ϕ, the ω' (parametrizing the correlated $\pi\rho$ exchange) and a fourth pole (denoted S). In what follows, we will assume that from these only the ϕ and the S couple to strangeness. Notice that it has recently been shown that there is no enhancement close to threshold of the isoscalar spectral function due to pion loops [20]. Furthermore, $\Lambda \sim 10\,\text{GeV}^2$ [7] separates the hadronic from the quark contributions, Q_0 is related to Λ_{QCD} and γ is the anomalous dimension,

$$F_i(t) \to (-t)^{-(i+1)}\left[\ln\left(\frac{-t}{Q_0^2}\right)\right]^{-\gamma}, \quad \gamma = 2 + \frac{4}{3\beta}\,, \quad i = 1, 2\,, \tag{4}$$

for $t \to -\infty$ and β is the one loop QCD β–function. In fact, the fits performed in [7] are rather insensitive to the explicit form of the asymptotic form of the spectral functions. To be specific, the additional factor $L(t)$ in Eq.(3) contributes to the spectral functions for $t > \Lambda^2$, i.e. in some sense parametrizes the intermediate states in the QCD regime, above the region of the vector mesons. The particular logarithmic form has been chosen for convenience. Obviously, the asymptotic behaviour is obtained by choosing the residues of the vector meson pole terms such that the leading terms in the $1/t$–expansion cancel. In practice, the additional logarithmic factor is of minor importance for the fit to the existing data. The number of isoscalar and isovector poles in Eq.(3) is determined by the stability criterion discussed in detail in [5][7]. In short, we take the minimum number of poles necessary to fit the data. Specifically, we have four isoscalar and three isovector poles. This fourth isoscalar pole is necessary since most isoscalar couplings are fixed (as described above) and otherwise we would not be able to fulfill the various normalization and superconvergence relations. We are left with three fit parameters, these are the masses of the third isovector and the fourth isoscalar pole as well as the residuum a_1^{ω}.

The spectral functions of the isoscalar form factors $F_{1,2}^{(0)}$ encode information about the strange vector current since the photon couples to a certain extent via mesons with strangeness (here the ϕ and the S) to the nucleon. Assuming that the strange form factors have the same large momentum fall–off as the isoscalar electromagnetic ones [4][6] and neglecting the small $\omega - \phi$ mixing, it is straightforward to extract the strange Dirac and Pauli form factors following the formalism outlined in [4][6]

$$F_1^s(t) = t\, L(t)\, a_1^\phi L_\phi^{-1} \frac{M_\phi^2 - M_S^2}{(t - M_\phi^2)(t - M_S^2)}\,,$$
$$F_2^s(t) = L(t)\, a_2^\phi L_\phi^{-1} \frac{M_\phi^2 - M_S^2}{(t - M_\phi^2)(t - M_S^2)}\,, \tag{5}$$

with $L_\phi^{-1} = 1/L(M_\phi^2)$. Clearly, the size of these strange form factors is given by the strength of the ϕ–nucleon couplings (as encoded in the residua $a_{1,2}^\phi$). In particular, we notice that the sign of the strange radius $r_{1,s}^2$ is determined from the sign of a_1^ϕ whereas the sign of the strange magnetic moment, $\mu_s = F_2^{(s)}(0)$, is fixed by the sign of the tensor coupling $\sim a_2^\phi$.

A best fit to the available data as compiled in [21] is obtained with $M_{S'} = 1.63\,\mathrm{GeV}$, $M_{\rho'''} = 1.72\,\mathrm{GeV}$ and $a_1^\omega = 0.677$ (for $g_{\phi NN} = -0.24$ and $\kappa_\phi = 0.2$). The χ^2/datum of the fit is 1.02. All contraints are fulfilled to high numerical accuracy. A detailed account of these results is given in [22]. The corresponding strange form factors are shown in fig. 3. Notice that $F_1^{(s)}(t)$ varies very weakly between $t = -1 \ldots - 10\,\mathrm{GeV}^2$. Furthermore, the strange magnetic moment and radius are $\mu_s = 0.003\,\mathrm{n.m.}$ and $r_s^2 = 0.002\,\mathrm{fm}^2$, respectively. These are orders of magnitude smaller than in previous analysis [4][6] where the ϕ pole subsumed the non–strange physics of the isoscalar spectral function in the mass region of about 1 GeV, i.e. the sizeable effect of the $\pi\rho$ correlations. The use of a more realistic spectral function based on the constraints from the NN interaction indeed leads to a reduction of the strange matrix elements as anticipated from the OZI rule. Of course, the analysis presented here can be sharpened by including the effects of $\omega - \phi$ mixing and by studying the dependence on the large-t behaviour of the strange form factors. This will, however, not change the main conclusion of our work, namely that the inclusion of correlated $\pi\rho$–exchange in the isoscalar spectral function for the nucleon electromagnetic form factors leads to a sizeable reduction of the strange vector current matrix elements in the proton. The experimental information concerning the strange form factors is thus eagerly awaited for.

Acknowledgements

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Figure Captions

Fig.1 Correlated $\pi\rho$ exchange missing in the Bonn potential.

Fig.2 Hadronic model for the ϕNN vertex consisting of Born terms and interaction diagrams (as indicated by the blobs).

Fig.3 Strange form factors $F_1^{(s)}(t)$ (solid line) and $F_2^{(s)}(t)$ (dashed line).

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Tables

	g	f
$\bar{K}K$	-0.32	-0.15
$\bar{K}^*K - \bar{K}K^*$	-0.39	-0.20
$\bar{K}^*K^*$	$+0.25$	$+0.40$
all mesons	-0.47	$+0.14$
$\Sigma\bar{\Sigma} + \Lambda\bar{\Lambda}$	$+0.23$	-0.19
Sum	-0.24	-0.05

Table 1: Various contributions to the ϕNN vector (g) and tensor coupling (f) (compare fig. 2). Given is also the sum of the individual contributions. In the case of the tensor coupling, the meson contributions can not simply be added since the spectral functions have different signs and the corresponding monopole fit to the sum of the spectral functions has a considerably harder form factor than the individual contributions. In contrast, for the vector coupling the monopole cut offs are all of comparable size and thus one can essentially add the individual contributions.

Figures

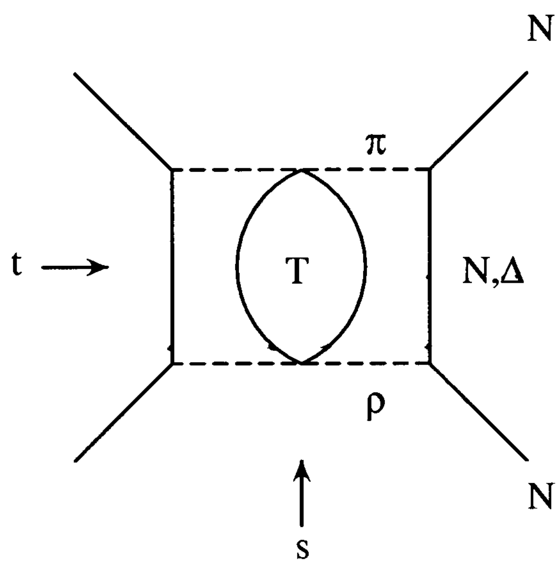

Figure 1

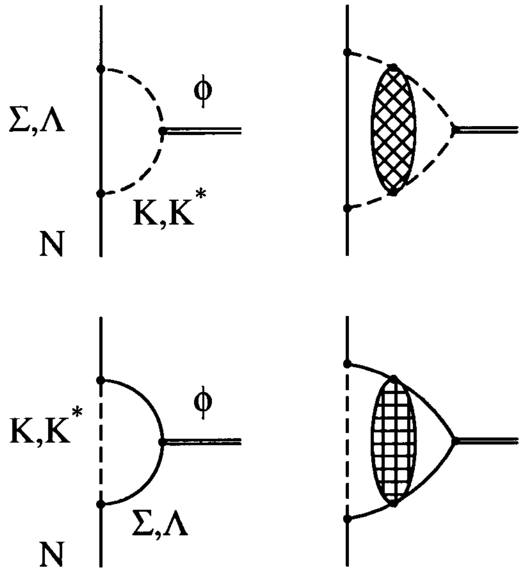

Figure 2

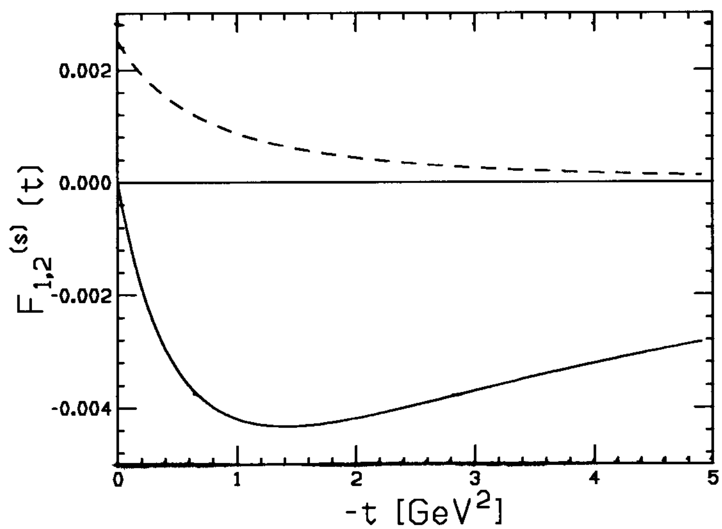

Figure 3