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HADRON RADII AND THE ORIGIN OF CHIRAL SINGULARITIES*

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Hadron radii are discussed with particular emphasis on chiral singularities m_{PS}^{-1} and log m_{PS}; those result from the photon coupling to the $(PS + PS)$ continuum and are, therefore, most visible in *isovector* electromagnetic radii. It is shown that OZI rule violation leads to m_K^{-1} and log m_K-terms in the *isoscalar* radii as well. The $\gamma\pi\pi\pi$ vertex function is free of chiral singularities and so are all quark core contributions. The $\Delta(1232)$ contribution is particularly important for the subleading log m_π-terms, where a delicate cancellation occurs with pion loop contributions of the same order.

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1. Introduction

There are more than 260 experimentally well-established particles known as "hadrons". For only a few their "size" has been determined experimentally. The "size" of any object depends on the nature of the (point-like) probe which is used in the scattering process. The systematics of hadron radii can, at least at a qualitative level, be obtained from an analysis of total hadron–proton cross sections, where the introduction of "effective" radii allows for a quasi-geometrical picture, see [1]. The emerging pattern for $p, \pi, K, \phi, J/\psi, \Lambda, \sum, \Xi, \overline{p}$ is strikingly similar to the one obtained from electromagnetic probes and in some cases $(\phi, J/\psi, \Lambda, \sum, \Xi, \overline{p})$ is the only empirical information on the hadron's size. A more accurate way[1] of probing

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[1] Pointlike particles (like the electron) have an accurately known QED vertex (in the one-photon approximation); this is not the case for the hadronic QCD vertices (the one-gluon approximation doesn't make sense here) involved in hadron–hadron collisions.

a hadron's size is by scattering elastically **point-like** particles off the extended hadron (or as in the case of the neutron $n(939)$ thermal neutrons have been scattered off atomically bound electrons). The fact that empirically $r_{e,\mu,\tau} \lesssim 10^{-3}\,\mathrm{fm}$ [2] points naturally at leptonic currents (neutral leptons are believed to be point-like, subject to (extremely difficult) precision tests of the weak interaction theory within the Standard Model). This simple fact has initially been widely used by Rutherford and Hofstätter in their historic experiments to obtain first information on nuclear and nucleonic radii (a hadron with its numerous different flavour "excitations" could be seen as the smallest nucleus including all its "isotopes"). A scattering process using the weak current instead would allow to determine a different size parameter, the weak hadronic size, and so on. On the theoretical side the size parameter ("rms radius") of a specific hadron is a distinctive property of that hadron's "form factor" which describes the current expectation value between the incoming and outgoing hadron. The simplest example is the electromagnetic form factor of the (pseudoscalar) lightest hadron, the pion $\pi\,(140)$, defined by

$$2m_\pi \langle \pi(p')|j^\mu(0)|\pi(p)\rangle = e_\pi (p'+p)^\mu\, F_\pi(q^2)\,, \tag{1}$$

where $q^\mu = (p'-p)^\mu$ and e_π is the pion's charge. The pion radius is now defined by[2] (with an obvious generalization for other hadrons)

$$\langle r_\pi^2\rangle = -\,6\frac{d}{dq^2}\,\frac{F_\pi(q^2)}{F_\pi(0)}|_{q^2\to 0^-}\,. \tag{2}$$

This definition is motivated by the (non-relativistic) definition of a rms radius $\langle r^2\rangle = \int d^3\vec{r}\, r^2\rho(r)/\int d^3\vec{r}\rho(r)$. The connection between these two definitions is the well-known fact that the form factor $F(q^2)$ (for example as in (1)) is the Fourier-transform of the corresponding (electric,magnetic (as in (1,2)), or axial charge, scalar, pseudoscalar, *etc.*) density $\rho(r)$ which in the non-relativistic limit is $\rho(r) = \psi^+(\vec{r})\psi(\vec{r}) = \overline{\psi(\vec{r})}\psi(\vec{r})$. In the fully relativistic situation (as is the case for the pion) the radius is **defined** by (2). The form factor is deduced directly from the measured differential cross sections. For electron-nucleon scattering two form factors $G_{E,M}^{(p,n)}(q^2)$ have to

[2] In the case of a neutral hadron $F(0) = 0\,(n,\pi^0,K^0,etc.)$ and $\langle r_{h^0}^2\rangle = -6\frac{d}{dq^2}\,F(q^2)|_{q^2=0}$ instead; for the neutron this leads to a negative $\langle r^2\rangle$. The sign is due to an inner positive core (p) and a negative cloud (π^-). Due to C-invariance the photon can not couple to the π^0 which is its own anti-particle, in line with a $(u\overline{u}, d\overline{d})$ quark content. Therefore, the π^0 has no electromagnetic form factor, unlike the $K^0 = d\overline{s}$ which is different from $\overline{K^0} = \overline{d}s$. Note that there is no contradiction to the existence of the electromagnetic polarizability for the π^0 which requires a Compton amplitude (2 photons in a contact interaction with the π^0, called a "seagull" diagram).

be taken into account. These can be disentangled from cross sectional data using the "Rosenbluth plot". They give rise, via a generalised Eq. (2), to different sizes $< r^2_{E,M}\rangle_{(p,n)}$ of proton (p) or neutron (n). In an analogous fashion a scalar radius of the pion and nucleon can be defined, although not directly accessible to experiment. A typical result is $< r_s^2\rangle = (0.54, 1.5)\,\mathrm{fm}^2$ for the (pion, nucleon) [3].

On the theoretical side hadron radii, like other low-energy parameters of hadrons, have entered "low-energy" theorems and other relations among low-energy parameters which have to be fulfilled by **any** low-energy approximation to proper QCD which carries the essential chiral symmetry of QCD, breaks it spontaneously in the same way as QCD and produces dynamically a non-vanishing quark scalar condensate and a "constituent" quark mass that satisfies the Goldberger–Treiman relation, see [4, 5]. The systematic low-energy expansion of QCD has become known under the label "χPTh". A sizeable number of low-energy parameters of χPTh depend on hadron radii. Furthermore, electromagnetic radii strongly influence the (usually well measured) electromagnetic mass shift of hadrons, see for example [5] and further references to earlier papers there. A further peculiarity is that chiral symmetry is realized in a different way for mesons and baryons. For mesons (notably the lightest meson, the pion) the chiral limit is $m_h \to 0$. While this is suggested by Goldstone's theorem (which implies that quark masses $m_q \to 0$ as well) and confirmed by the near masslessness of the pion, the chiral limit for baryons is far from undisputed. There are strong indications that the chiral limit in the baryon sector would correspond to the heavy fermion limit [6] (which implies that the lightest baryon, the nucleon, would be furthest away from this limit, unlike the pion which is closest amoung all mesons to the chiral limit). Hadronic radii have recently been studied for hadrons which are embedded in a hot and/or dense medium instead of in a vacuum [7] (if not otherwise stated values for radii will refer to the vacuum). This leads to the notion of the "swelling" of nucleons in a medium. We will briefly return to this issue below.

2. Chiral singularities

Chiral symmetry and its observed breaking pattern is reflected in hadron radii in the form of $1/m_{ps}$ and log (m_{ps}) "chiral singularities". This was first noted more than 25 years ago [8],

$$
\begin{aligned}
\langle r^2_{\pi^\pm}\rangle &= -\frac{2}{(4\pi f_\pi)^2}\log\left(\frac{m_\pi}{\mu_\pi}\right) + \text{f.t.}\,,\\
\langle r_1^2\rangle^{I=1} &= -\frac{1+5g_A^2}{8\pi^2 f_\pi^2}\log\left(\frac{m_\pi}{\mu_N}\right) + \text{f.t.}\,,
\end{aligned}
$$

$$\left(\frac{\mu_p-\mu_n}{\mu_N}-1\right)\langle r_2^2\rangle^{I=1} \;=\; \left(\frac{g_{\pi NN}^2}{8\pi m_N^2}\right)\frac{m_N}{m_\pi}+\left(\frac{3g_{\pi NN}^2}{2\pi^2 m_N^2}\right.$$
$$\left.+\frac{1}{4\pi^2}\left[\frac{\partial}{\partial\nu}A^{(-)}(\nu,0)\right]_{\nu=0}\right)\log\frac{m_\pi}{\mu_N}+\text{f.t.} \quad (3)$$

The origin of these chiral singularities is in all the above cases the $\pi\pi$ intermediate state (or in general the $(PS+PS)$ state) see Figs. 1, 2. Empirically $\langle r^2_{\pi^\pm}\rangle = [0.657(12)\text{fm}]^2$ (Ref. [9]), $\langle r^2_{K^\pm}\rangle = [0.53(5)\text{fm}]^2$ (Ref. [10a]), $\langle r^2_{K^0}\rangle = -0.054(26)\,\text{fm}^2$ (Ref. [10b]), $\langle r^2_E\rangle_p = (0.838(12)\text{fm})^2$, (Ref. [11]) or $(0.862(12)\text{fm})^2$, (Ref. [12]), $\langle r^2_E\rangle_n = -0.113\pm 0.003\pm 0.004\,\text{fm}^2$ (Ref. [13]), $\langle r^2_M\rangle_p = (0.853(9)\,\text{fm})^2$, (Ref. [14]), $\langle r^2_M\rangle_n = (0.889(9)\text{fm})^2$, (Ref. [14]) or $(0.85(1)\text{fm})^2$ (Ref. [11]), and the negative sign of $\langle r^2_{K^0}\rangle$ hints at a positively charged core $K^*(892)^+$ and a negatively charged pion cloud. This is in line with the non-relativistic quark model which has the heavier $\overline{s}$, which is *positively* charged, closer to the origin than the *negative* but much lighter d-quark. Note the unspecified scale parameters $\mu_{\pi,N}$ in Eq. (3). In the

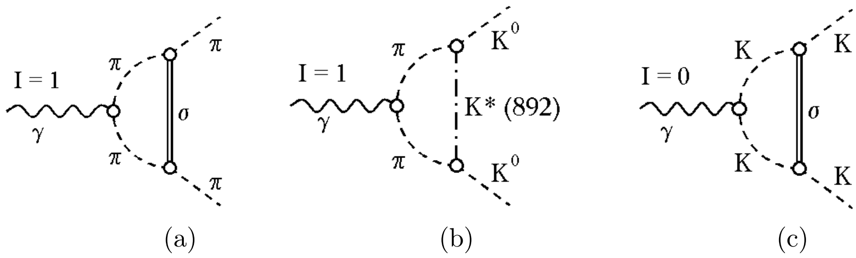

Fig. 1. $(a+b)$ Pion cloud corrections to pion (a) and K^0 electromagnetic form factor; (c) contribution which gives rise to a log m_K-term in $K^\pm$ electromagnetic form factor. The circles respresent vertex functions due to sub-hadronic degrees of freedom.

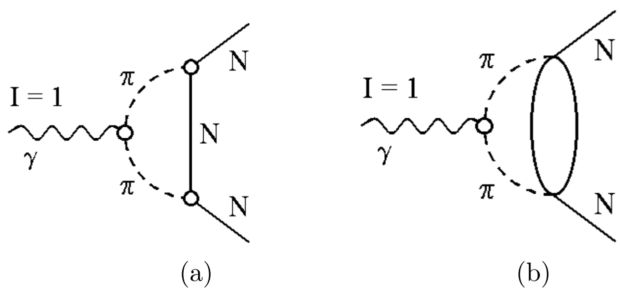

Fig. 2. Isovector electromagnetic form factor of the nucleon. (a) Born term, (b) full πN scattering amplitude.

simplest pion model where only a σ-meson is exchanged between the two pions in the intermediate state (Fig. 1(a)) $\mu_\pi = m_\sigma$ results [15]. In the simplest model of nucleon mesonic substructure only a $g.s.$ nucleon is exchanged between the two pions in the intermediate state (Fig. 2(a)) and $\mu_\pi = m_N$ with

$$g_A \equiv 1, \frac{\partial}{\partial \nu} A^{(-)}(\nu, 0)|_{\nu=0} \ \equiv \ 0 \tag{4}$$

results [16]. This is perhaps not surprising as the $\pi N \to \pi N$ amplitude enters only at the tree level (Fig. 2(a)); if, however, more than the Born term is included (Fig. 2(b)) one encounters at least two-loop contributions and g_A and $\frac{\partial}{\partial \nu} A^{(-)}(\nu, 0)|_{\nu=0}$ will start to deviate from 1 and 0, respectively. The fact that g_A remains unrenormalized at one-meson loop level [16] has been confirmed at the quark level in a different chiral model (the $0(1/N_c)$ corrections to the Nambu–Jona–Lasinio model, see Ref. [4]).

Of course the findings (3) are confirmed by χPTh; a notable exception are only Skyrme-type models with pion masses [17–19], which find instead of (3),

$$\langle r_1^2 \rangle^{I=1}_{\text{Skyrme}} \sim m_\pi^{-1} \,. \tag{5}$$

Unfortunately, no results for $\langle r_2^2 \rangle^{I=1}_{\text{Skyrme}}$ have been reported which could be compared with (3).

The reason for the difference between (5) and (3) as suggested in [18] is the non-commutativity of the chiral limit $m_\pi \to 0$ with the large-N_c limit. The Skyrme model takes the limit $N_c \to \infty$ implicitly from the beginning while the relations (3) are a finite- N_c result. Ref. [17] offers a mechanism by which the discrepancy would disappear for $N_c \to \infty$: the $B = \Delta(1232)$-exchange contribution to the relevant diagram (see Fig. 1(b) of [16])

$$\gamma\, N \to \gamma\, \pi\, B \to \pi\, B \to N \tag{6}$$

cancels exactly the chiral log-term in (3) for $N_c \to \infty$, and results in a form (5). It is important to note that the m_π^{-1} singularity in (3) is a result of the fermionic nature of the exchanged particle (baryon B) in (6). Hence there is only a log m_π-term in the corresponding mesonic isovector radius, see first equation of (3), which results from the $\sigma-$meson exchange[3]. The spin- 3/2 nature of the $\Delta(1232)-$exchange could explain the removal of the exact cancellation of the coefficient of m_π^{-1} due to N- exchange in $\langle r_1^2 \rangle^{I=1}$. It

[3] The meson propagator $\sim \frac{1}{m^2 - p^2}$ in loop integrals $\int d^4p...$ is replaced with a fermion propagator $\sim \frac{\not{p}+...}{m^2-p^2}$ in loop integrals $\int d^4p...$ which results in a higher singularity due to $\not{p}$ in the numerator.

is important to further clarify this point in the soliton as well as the $N_c < \infty$ sector. The $\Delta(1232)$ contribution plays a similar role in the electromagnetic polarizabilities of the proton (for which there are recent, rather accurate experimental values [20]) which are proportional to m_π^{-1} and do not have a log m_π-term, due to a peculiar cancellation of $\Delta(1232)$ tree and pion loop Compton amplitudes, see [21].

An independent way of studying the origin of the chiral singularities is via dispersion relations. If the form factors $G_{E,M}^{(S,V)}(t)$ fulfill an unsubtracted dispersion relation[4] then

$$\langle r^2 \rangle = \frac{6}{\pi} \int\limits_{t_o}^{\infty} dt' \frac{ImG(t')}{t'^2} \tag{7}$$

with $t_0 = 4m_\pi^2 (9m_\pi^2)$ for $I = 1(0)$. The chiral limit is obtained for $t_0 \to 0$ (*i.e.* t' runs between 0 and ∞). One finds $\frac{ImG_{E,M}^V(t)}{t^2} \sim \frac{1}{t^2}, \frac{1}{t}$ for $t \to 0$ resulting in m_π^{-1} and log m_π-terms in $\langle r^2 \rangle^{I=1}$. Due to a delicate cancellation which occurs only for N-exchange contributions, the m_π^{-1}-term in $\langle r_1^2 \rangle^{I=1}$ vanishes while both terms survive in the combination $\langle r_2^2 \rangle^{I=1}$. With respect to the *isoscalar* form factors one finds instead $\frac{ImG_{E,M}^S(t)}{t^2} \sim t, \sqrt{t}$ for $t \to 0$. Hence there are only finite terms in $\langle r_{1,2}^2 \rangle^{I=0}$ for $m_\pi \to 0$.

A striking feature of (3) is that all chirally singular radii are **isovector**, none of them is **isoscalar** in nature. That fact is quite astonishing as the probing virtual photon is known to be isovector (ρ-like) **and** isoscalar ($\omega-, \phi-, J/\psi-$like) with comparable weights. Therefore, it is the ρ-like part of the photon and its strong coupling to the 2π-intermediate state[5] which causes the chiral singularity in the isovector radii. Knowing that the ω-like part of the photon couples to the $\pi\rho = 3\pi$ intermediate state (and only very little to the 2π-state via $\rho - \omega$ mixing) it is not surprising that the $\pi\rho = 3\pi$ intermediate state does not generate any chiral singularities, hence the isoscalar radii are free from singularities from this intermediate state. In the vector-meson-dominance picture, part of the photon "behaves" like a virtual $\phi(1019)$ which then naturally couples to the $K\overline{K}$ intermediate state. Following the same arguments which led to Eq. (3), it is not difficult to see that for $m_K \to 0$ the dominant contributions to the isoscalar $\langle r_{1,2}^2 \rangle^{I=0}$ will be $\log(\frac{m_K}{\mu_N})-$ and m_K^{-1}- terms at the scale $\mu_N = m_\Lambda = 1.12\,\mathrm{GeV}$, because now the lightest exchanged baryon will be $\Lambda(1115)$ or $\sum^{\pm}(1197)$, see Fig. 3. If

[4] The existence of dispersion relations in proper QCD is not undisputed, see Ref. [22].
[5] The virtual ρ couples to the 2π continuum even below the ρ-resonance, close to threshold, see discussion in [23, 24].

one draws Fig. 3 as a quark-line diagram, the strange quark line is closed and "disconnected" from the non-strange quark lines in the nucleon; according to the Okubo–Zweig–Iizuka (OZI) rule such diagrams should be supressed (here manifest in a small ΛKN and $\sum KN$ coupling). Strong violations of the OZI rule have, however, been discussed, also in connection with the strange vector form factors of the nucleon, see discussion in refs [25, 26].

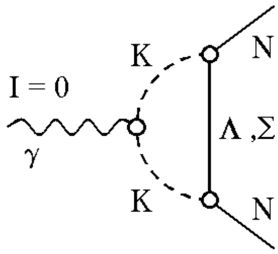

Fig. 3. Isoscalar electromagnetic form factor contribution due to the Kaon cloud.

It might seem unphysical to consider the limit $m_K \to 0$ for the massive $K(494)$. In most $\mathrm{SU}(6)_{\mathrm{flavour}}$ models, however, one separates quarks into three "light" quarks (u, d, s) and three heavy quarks (c, b, t). The former set would lend itself naturally to the chiral limit $m_q \to 0$ (mesonic chiral limit), whereas the latter part seems closer to the heavy quark limit $m_Q >> m_q$ $(Q = c, b, t$ and $q = u, d, s)$ *i.e.* the baryonic chiral limit. The singularities in baryon radii occur for $m_{\pi,K} \to 0$ and might not be directly related to the baryonic chiral limit.

According to the relevant diagram Fig. 1(b) we expect a $\log \frac{m_\pi}{\mu_K}$–term with $\mu_K = m_{K^*} = 892\,\mathrm{MeV}$ in $\langle r^2_{K^0} \rangle$ as well. The absence of Fig. 1(b) for the charged K leads to a $\log \frac{m_K}{m_\sigma}$–term instead, due to Fig. 1(c).

We note in passing that one can obtain **analytical** results for $\langle r^2_\pi \rangle$ [15] and $\langle r^2_{1,2} \rangle^{I=0,1}$ [16]. Without any prejudice towards the correct answer to the question "do hadrons swell when embedded in a medium?" we would like to mention that those analytical results for radii allow us to study these hadronic radii as a function of the parameters $m_{\pi,\sigma}, m_N, g_{\pi qq}, g_{\pi NN}, g_{\sigma NN}, f_\pi$. Medium effects might change those parameters and, hence, could influence $\langle r^2 \rangle$ in a well-defined way, within the approximations of [15] and [16]. The analysis of [16] clearly showed that the σ-meson cloud around the nucleon is important to reproduce the nucleon radii, and the anticipated swelling of the nucleon in a medium could conceivably come from an enhancement of the σ-meson cloud due to medium effects. We find that changing the σ-meson mass and coupling has a more dramatic effect than changing the pion mass and coupling[6]. An anticipated "swelling of the nucleon" could also come from a medium-enhanced three-pion state contribution (the $\gamma\,\pi\,\pi\,\pi$ anomaly); in that case the medium effects should conspire to increase the parameter

$$\kappa = \frac{9 g^3_{\pi NN}\, g_{\rho\pi\pi}\, M_N\, g_{\omega\pi\rho}}{128\pi^4\, g_\omega}$$

(see discussion in [16] and [27]).

3. Quark core of nucleon and pion

We have recently extended the investigation in [16] to include (quark) substructure of π, σ and N in a fully gauge-invariant manner [28]. Without (quark) substructure the nucleon electromagnetic radii generally come out too small, although their magnetic moments can be reproduced. The simplest extension of the sigma model with a fixed cut-off is the replacement of the point couplings πNN, σNN, $\gamma\pi\pi$, γNN, with vertex functions which naturally cut down large momentum transfers due to (quark or higher-order mesonic) substructure of the hadrons involved. The $\gamma\pi\pi$ vertex function, including the all-important non-perturbative $q\overline{q}$ substructure of the (Goldstone) pion will be taken from the NJL model calculation of Ref. [15]. For other vertex functions we use various relativistic quark model calculations [27, 29]

Not surprisingly we find no factorization of form factors due to the quark core and those due to the meson cloud of the nucleon. Therefore, without additional assumptions, in general

$$\langle r^2\rangle_N \neq \langle r^2\rangle_{\rm core} + \langle r^2\rangle_{\rm cloud}$$

independent of the specific form of the quark form factors. Quite generally one finds that chiral singularities do **not** come from the quark core of either π, σ or N. Nevertheless, the quark core of hadrons is important to reproduce the empirical radii. To illustrate this we quote the following results for the singular terms alone, with $m_\pi = 140\,\mathrm{MeV}$, and compare with the empirical values,

$$\langle r_1^2\rangle^{I=1} = \begin{cases} 0.96\,\mathrm{fm}^2 & \text{for } g_A = 1.25 \\ 0.65\,\mathrm{fm}^2 & \text{for } g_A = 1 \\ 0.58\,\mathrm{fm}^2 & \text{empirically} \end{cases}$$

[6] In models where the $\sigma-$meson couples to the isoscalar-scalar 2π continuum, only this ($I = 0$, $J = 0$) part of the 2π continuum would possibly be enhanced (*i.e.* not Fig. 2, but Fig. 1(c) of [16] is responsible for the effect).

$$\langle r_2^2\rangle^{I=1} = \begin{cases} \leq 0 & \text{for } \frac{\partial}{\partial\nu}A^{(-)}(\nu,0)|_{\nu=0} \equiv 0 \\ 0.77\,\mathrm{fm}^2 & \text{empirically} \end{cases}$$

Without the finite terms $\langle r_2^2\rangle^{I=1}$ would be negative for $m_\pi = 140\,\mathrm{MeV}$ (the positive $\frac{1}{m_\pi}$-term and the negative $\log m_\pi$-term compete in $\langle r_2^2\rangle^{I=1}$ with the $\frac{1}{m_\pi}$-term taking over only for relatively small values of m_π).

In order to study the effects from the quark core and the meson cloud in a meaningfull way one has to compare with all available form factor data. Some form factors turn out to be more sensitive to the cloud, others are dominated by the quark core. It is useful to compare with phenomenological fits which are available for most form factors. They reproduce at least the normalization $G(0)$ and slope at $q^2 = 0$ (*i.e.* the radius) in the space-like region of that form factor $G(q^2) = G(0)[1 + \frac{1}{6}q^2\langle r^2\rangle + O(q^4)]$. In particular, for $q^2 \leq 0$

$$\begin{aligned} G_E^p(q^2) &= [1 - q^2\sqrt{2}\,\mathrm{GeV}^{-2}]^{-2} = G_D(q^2)\,, \\ G_M^p(q^2) &= \mu_p G_E^p(q^2)\,, \\ G_E^n(q^2) &= -\frac{q^2}{q_0^2}\,\mathrm{e}^{q^2 R_0^2},\ G_M^n(q^2) = \mu_n G_D(q^2) \end{aligned}$$

with $\mu_{p,n} = (2.793, -1.913)$, $q_0^2 = 1.95\,\mathrm{GeV}^2$, $R_0^2 = 3.8\,\mathrm{GeV}^{-2}$, and

$$G_A(q^2) = g_A \begin{cases} [1 - q^2/M_A^2]^{-2} & \text{see Ref. [19]} \\ \left[1 - q^2/2M_\rho^2\right]^{-1} \exp\left[\frac{1}{6}q^2 R_{\mathrm{L}}^2/(1 - \frac{q^2}{4M_N^2})\right], & \text{Ref. [30]} \\ \exp\left[q^2/\Lambda^2\right] & \text{see Ref. [31]} \end{cases}$$

$$F_\pi(q^2) = [1 - q^2/M_\rho^2]^{-1}$$

with $M_\rho = 770\,\mathrm{MeV}$, $M_A = (900-1100)\,\mathrm{MeV}$, $R_{\mathrm{L}}^2 = 6\,\mathrm{GeV}^{-2}$, $\Lambda = 770\,\mathrm{MeV}$. The corresponding charge distributions $\rho(\vec{r}) = \int \frac{d^3q}{(2\pi)^3}\mathrm{e}^{i\vec{q}\cdot\vec{r}}G(\vec{q}^2)$ can be found in Fig. 3 of Ref. [23], where the Breit-frame ($q_0 = 0, q^2 = -\vec{q}^2$) has been used.

Other form factors are related to those above

$$G_{\pi NN}(q^2) = \frac{M_N}{f_\pi}[G_A(q^2) - \eta G_p^{\mathrm{quark}}(q^2)]. \tag{8}$$

The pseudoscalar form factor $G_p(q^2) = G_p^{\mathrm{quark}}(q^2) + G_p^\pi(q^2)$, with $G_p^{\mathrm{quark}}(q^2)$ in the Goldberger–Treiman relation (8), and $G_p^\pi(q^2) = 4M_N F_\pi \frac{G_{\pi NN(q^2)}}{m_\pi^2 - q^2}$, the

PCAC relation. Parametrizations for time-like q^2 can be found in Ref. [23]. The last relations indicate where a separation into quark core and meson cloud in form factors might make sense. The non-renormalization of g_A by one-meson loop contributions [15,16] implies that $G_A(q^2)$ should (at least for moderate q^2) be dominated by the quark core of the nucleon. The size parameter, $\langle r_A^2 \rangle^{\frac{1}{2}} = (0.5-0.6)$ fm, comes out significantly smaller than the electromagnetic nucleon radii and hints at a non-negligible extent of the quark core. For a proper, fully relativistic description of the low-energy parameters of the pion, the confinement aspect of QCD seems to be irrelevant [4, 5, 15]. This is not so for the nucleon where a systematic investigation of possible quark-confining potentials is called for. For free nucleons and mesons (excluding the Goldstone pion) bag-type Lorentz-scalar confining potentials with additional Lorentz-vector components (due to the residual one-gluon exchange (OGE) interaction) have been widely discussed. Effects of Lorentz-vector quark confinement in extensions of the NJL model for mesons have been discussed [32]. In the baryon sector a peculiarity of Lorentz-vector confinement of Dirac- particles in the form of the "Klein paradox" has to be taken into account : spin-$\frac{1}{2}$ particles (quarks) can not be confined by a *dominant* Lorentz-vector confining potential. For an arbitrary scalar confinement there exists, therefore, an upper limit on the strength of a vector confinement potential, so that the combination of both still allows for quark bound states. This is unlike the situation in the $q\overline{q}$ sector which does not constrain the strength of the Lorentz-vector component. We note here also, that a sizeable Lorentz-vector potential in the Dirac-equation for the 3 quarks is necessary to avoid a large spin-orbit force in the baryon spectrum. For a combination of Lorentz-scalar confinement $M(r) = c_2 r^2$ (with $c_2 = 1.0239\,\text{GeV}\,\text{fm}^{-2}$ so that the quark core size is 0.6 fm [33]) and Lorentz-vector confinement

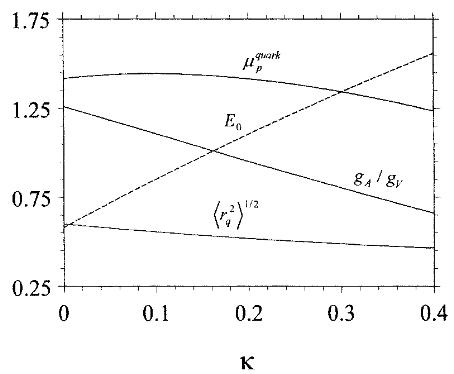

Fig. 4. $g_A/g_V, \mu_p^{\text{quark}}$ in units of $\mu_N, \langle r_q^2 \rangle^{\frac{1}{2}}$ [fm] and E_0 [GeV], the quark eigen energy, as a function of κ. Here we have used $M(r) = c_2 r^2$, see text.

$V(r) = \alpha_s \ln(r/r_0)$ or $V(r) = \kappa r \mathrm{e}^{-\mu r}$ we have obtained g_A/g_V, $\langle r_q^2 \rangle^{\frac{1}{2}}$, E_0, and μ_p^{quark} as a function of κ in Fig. 4 (a similar tendency is found for α_s variations, see Ref. [33]). A full calculation of nucleon form factors (as in [28]), including the quark core, will limit the range of possible parameters κ or α_s in $V(r)$ and c_n in $M(r) = c_n r^n$. Again, medium effects will have an influence on κ, α_s and c_n, on one hand, and on the meson cloud, on the other hand.

4. Outlook

We plan to include the kaon cloud in our new approach [16, 28] and study m_K^{-1}, log m_κ-terms in the isoscalar nucleon radii. The coefficients of those terms depend on OZI-rule violating couplings $g_{\Lambda KN}$ and $g_{\Sigma KN}$. The strangeness content of the nucleon will be influenced by the magnitude of these couplings. We hope to report soon on our findings.

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