Sigma-term physics in the perturbative chiral quark model

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We apply the perturbative chiral quark model at one loop to analyze meson-baryon sigma terms. Analytic expressions for these quantities are obtained in terms of fundamental parameters of low-energy pion-nucleon physics (weak pion decay constant, axial nucleon coupling, strong pion-nucleon form factor) and of only one model parameter (radius of the nucleonic three-quark core). Our result for the πN sigma term $\sigma_{\pi N} \approx 45$ MeV is in good agreement with the value deduced by Gasser, Leutwyler, and Sainio using dispersion-relation techniques and exploiting the chiral symmetry constraints.

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I. INTRODUCTION

The meson-nucleon sigma terms are fundamental parameters of low-energy hadron physics since they provide a direct measure of the scalar quark condensates in baryons and constitute a test for the mechanism of chiral symmetry breaking (for a review and recent advances in sigma-term physics see Refs. [1-4]). In particular, sigma terms pose an important test for effective quark models in the low-energy hadron sector, since these quantities are mostly determined by the quark-antiquark sea and not by the valence quark contribution. The problem of the meson-nucleon sigma terms is also closely related to properties of light hadron phenomenology such as the chiral expansion of baryon masses, the pseudoscalar meson-nucleon scattering lengths [5-7], properties of hadronic atoms [8], and nuclear matter at finite temperature or density [9,10].

In current algebra the nucleon sigma term σ_{NN}^{ba} is derived from the low-energy theorem for meson-nucleon scattering. It is given by the double commutator of the axial-vector charge Q_5^a with the chiral symmetry breaking (χ SB) part of the Hamiltonian $\mathcal{H}_{\chi SB}$, averaged over nucleon states $|N\rangle$ [2]:

$$\sigma_{NN}^{ab} \doteq \langle N | [Q_5^a, [Q_5^b, \mathcal{H}_{\chi SB}]] | N \rangle.$$
(1)

In the QCD Hamiltonian the source of chiral symmetry breaking is simply the quark mass term with

$$\mathcal{H}_{\chi SB}^{QCD} = \bar{q} \mathcal{M} q = m_u \bar{u} u + m_d \bar{d} d + m_s \bar{s} s = \sum_{i=u,d,s} m_i S_i^{QCD},$$
(2)

where $\mathcal{M} = \text{diag}\{m_u, m_d, m_s\}$ is the mass matrix of current quarks and S_i^{QCD} is the scalar density operator corresponding to the quark of *i*th flavor. After straightforward use of current algebra one obtains

$$\sigma_{NN}^{ab} = \langle N | \left\{ \frac{\lambda^a}{2}, \left\{ \frac{\lambda^b}{2}, \mathcal{M} \right\} \right\} | N \rangle, \qquad (3)$$

where λ^a $(a=1,\ldots,8)$ are the corresponding Gell-Mann flavor matrices and curly brackets refer to the anticommutator. In this paper we perform the calculations of the πN , KN, and ηN sigma terms in the isospin symmetry limit with $m_u = m_d = \hat{m}$. With the use of key Eq. (3) the quantities of interest are given by

$$\sigma_{\pi N} \doteq \sigma_{pp}^{11} = \hat{m} \langle p | \bar{u}u + \bar{d}d | p \rangle, \qquad (4)$$

$$\sigma_{KN}^{u} \doteq \sigma_{pp}^{44} = \frac{\hat{m} + m_{s}}{2} \langle p | \bar{u}u + \bar{s}s | p \rangle,$$

$$\sigma_{KN}^{d} \doteq \sigma_{pp}^{66} = \frac{\hat{m} + m_{s}}{2} \langle p | \bar{d}d + \bar{s}s | p \rangle,$$

$$\sigma_{\eta N} \doteq \sigma_{pp}^{88} = \frac{1}{3} \langle p | \hat{m}(\bar{u}u + \bar{d}d) + 2m_{s}\bar{s}s | p \rangle,$$

where $|p\rangle$ denotes a one-proton state. By analogy we can define the sigma terms related to the Δ isobar by changing the proton state $|p\rangle$ to the corresponding state $|\Delta\rangle$ from the isomultiplet Δ (1230).

The pion-nucleon sigma term $\sigma_{\pi N}$ is equivalent to the value of the scalar nucleon form factor $\sigma(t)$ at zero value of momentum transfer squared: $\sigma_{\pi N} = \sigma(0)$ [7]. On the other hand, the pion-nucleon σ term is also related to the nucleon mass m_N by means of the Feynman-Hellmann (FH) theorem [1,11]

$$\sigma_{\pi N} = \hat{m} \frac{\partial m_N}{\partial \hat{m}}.$$
(5)

Both relations, $\sigma_{\pi N} = \sigma(0)$ and Eq. (5), are quite crucial as a consistent check of any approach applied to the study of the pion-nucleon σ term.

Other quantities which we consider are the strangeness content of the nucleon y_N and the isovector KN sigma term $\sigma_{KN}^{I=1}$:

$$y_{N} = \frac{2\langle p|\bar{s}s|p\rangle}{\langle p|\bar{u}u + \bar{d}d|p\rangle}, \quad \sigma_{KN}^{I=1} = \frac{\hat{m} + m_{s}}{4} \langle p|\bar{u}u - \bar{d}d|p\rangle,$$
(6)

where the strange quark condensate $\langle p|\bar{ss}|p\rangle$ is related to the nucleon mass by use of the FH theorem with $\langle p|\bar{ss}|p\rangle = \partial m_N / \partial m_s$. The isovector *KN* sigma term is related to the flavor-asymmetric condensate $\langle p|\bar{u}u - \bar{d}d|p\rangle$ of the proton. With the definitions of y_N and $\sigma_{KN}^{I=1}$ we can relate *KN* and ηN sigma terms to the πN sigma term as

$$\sigma_{KN}^{u} = \sigma_{\pi N} (1 + y_N) \frac{\hat{m} + m_s}{4\hat{m}} + \sigma_{KN}^{I=1}, \quad \sigma_{KN}^{d} = \sigma_{KN}^{u} - 2\sigma_{KN}^{I=1},$$
(7)

$$\sigma_{\eta N} = \sigma_{\pi N} \frac{\hat{m} + 2y_N m_s}{3\hat{m}}.$$

In the literature other definitions of the KN sigma terms can be found, namely Refs. [12-14],

$$\sigma_{KN}^{(1)} \equiv \sigma_{KN}^{u} \text{ and } \sigma_{KN}^{(2)} = \frac{\hat{m} + m_s}{2} \langle p | -\bar{u}u + 2\bar{d}d + \bar{s}s | p \rangle.$$
(8)

In the following we consider the relativistic quark model suggested in Ref. [15] and recently extended in Ref. [16] for the study of low-energy properties of baryons, which are described as bound states of valence quarks surrounded by a cloud of Goldstone bosons (π, K, η) as required by chiral symmetry. Similar models have been studied in Refs. [17– 19]. We refer to our approach as the perturbative chiral quark model (PCQM). The PCQM is based on an effective chiral Lagrangian describing quarks as relativistic fermions moving in a self-consistent field (static potential). The latter is described by a scalar potential S providing confinement of quarks and the time component of a vector potential $\gamma^0 V$ responsible for short-range fluctuations of the gluon field configurations [20]. Obviously, other possible Lorenz structures (e.g., pseudoscalar or axial) are excluded by symmetry principles. The model potential defines unperturbed wave functions of quarks which are subsequently used in the calculation of baryon properties. Interaction of quarks with Goldstone bosons is introduced on the basis of the nonlinear σ model [21]. When considering mesons fields as small fluctuations we restrict ourselves to the linear form of the mesonquark-antiquark interaction. With the derived interaction Lagrangian we do our perturbation theory in the expansion parameter 1/F (where F is the pion leptonic decay constant in the chiral limit). We also treat the mass term of the current quarks as a small perturbation. Dressing the baryon threequark core by the cloud of Goldstone mesons corresponds to inclusion of the sea-quark contribution. All calculations are performed at one loop or at order of accuracy $o(1/F^2, \hat{m}, m_s)$.¹ The chiral limit with $\hat{m}, m_s \rightarrow 0$ is well defined. Electromagnetic gauge invariance is guaranteed at all steps of the calculations in a specific frame, that is the Breit frame (for details, see Ref. [16]).

In the present paper we proceed as follows. First, we describe the basic notions of our approach: the underlying effective Lagrangian, the Dirac equation for quarks in the model potential, choice of free parameters, and fulfillment of low-energy theorems, and finally, key perturbative equations for the calculation of baryon masses and sigma terms. Other details (renormalization, gauge invariance, and application to electromagnetic properties of nucleons) can be found in Ref. [16]. Next, we concentrate on the detailed analysis of the meson-nucleon sigma terms in our approach: the chiral expansion of the πN sigma term in powers of pion mass and, after taking into account the kaon and η -meson loop contributions, model predictions for the KN, ηN sigma terms, and the strangeness content of the nucleon. Finally, we present our predictions for the Δ -isobar sigma terms. We compare all our predictions to results of other approaches, in particular, chiral perturbation theory (ChPT), heavy baryon chiral perturbation theory (HBChPT), and lattice QCD.

II. PERTURBATIVE CHIRAL QUARK MODEL

We start with the zeroth-order Lagrangian describing relativistic quarks in an effective static potential $V_{eff}(r) = S(r)$ + $\gamma^0 V(r)$ with $r = |\vec{x}|$ [15,18,22],

$$\mathcal{L}^{(0)}(x) = \overline{\psi}(x) [i\vartheta - \mathcal{M} - S(r) - \gamma^0 V(r)] \psi(x), \qquad (9)$$

which obviously is not chiral invariant due to the presence of the quark mass term \mathcal{M} and the scalar potential S(r). To recover chiral invariance (of course, in the limit $\mathcal{M} \rightarrow 0$) we introduce the interaction of quarks with the octet of Goldstone pseudoscalar mesons (π, K, η) by use of the nonlinear σ -model ansatz [21]. We define the chiral fields using the *exponentional parametrization* $U = \exp[i\hat{\Phi}/F]$, where *F* is the pion decay constant in the chiral limit. The octet matrix $\hat{\Phi}$ of pseudoscalar mesons is defined as

$$\frac{\hat{\Phi}}{\sqrt{2}} = \sum_{i=1}^{8} \frac{\Phi_i \lambda_i}{\sqrt{2}}$$

$$= \begin{pmatrix} \pi^{0}/\sqrt{2} + \eta/\sqrt{6} & \pi^+ & K^+ \\ \pi^- & -\pi^{0}/\sqrt{2} + \eta/\sqrt{6} & K^0 \\ K^- & \bar{K}^0 & -2\eta/\sqrt{6} \end{pmatrix}.$$
(10)

Kinetic term of meson fields \mathcal{L}_{Φ} and their interaction Lagrangian with quarks \mathcal{L}_{int} are given by

¹Throughout we use the Landau symbols O(x) [o(x)] for quantities that vanish like x [faster than x] when x tends to zero [8].

$$\mathcal{L}_{\Phi} = \frac{F^2}{4} \operatorname{Tr}[\partial_{\mu} U \partial^{\mu} U^{\dagger}] = \frac{1}{2} (\mathcal{D}_{\mu} \Phi_i)^2$$
(11)

and

$$\mathcal{L}_{int} = -\bar{\psi}(x)S(r) \left[\frac{U+U^{\dagger}}{2} + \gamma^5 \frac{U-U^{\dagger}}{2} \right] \psi(x)$$
$$= -\bar{\psi}(x)S(r) \exp\left[i\gamma^5 \frac{\Phi}{F} \right] \psi(x), \qquad (12)$$

where \mathcal{D}_{μ} is the covariant chiral derivative [17], defined as

$$\mathcal{D}_{\mu}\Phi_{i} = \partial_{\mu}\Phi_{i} + \left(F\sin\frac{\Phi}{F} - \Phi\right)\partial_{\mu}\left(\frac{\Phi_{i}}{\Phi}\right) \text{ with } \Phi = \sqrt{\Phi_{i}^{2}}.$$

The resulting Lagrangian $\mathcal{L}_{full} = \mathcal{L}_{inv} + \mathcal{L}_{\chi SB}$ contains a chiral-invariant piece \mathcal{L}_{inv}

$$\mathcal{L}_{inv}(x) = \overline{\psi}(x) [i\partial - \gamma^0 V(r)] \psi(x) + \mathcal{L}_{\Phi} + \mathcal{L}_{int}$$
(13)

and includes a term $\mathcal{L}_{\chi SB}$ which explicitly breaks chiral symmetry:

$$\mathcal{L}_{\chi SB}(x) = -\bar{\psi}(x)\mathcal{M}\psi(x) - \frac{B}{8}\mathrm{Tr}\{\Phi, \ \{\Phi, \ \mathcal{M}\}\}$$
(14)

corresponding to the mass terms for quarks and mesons. Here $B = -\langle 0 | \bar{u}u | 0 \rangle / F^2$ is the low-energy constant which measures the vacuum expectation value of the scalar quark densities in the chiral limit [23]. We rely on the standard picture of chiral symmetry breaking [23] and for the masses of pseudoscalar mesons we use the leading term in their chiral expansion (i.e., linear in the current quark mass)

$$M_{\pi}^2 = 2\hat{m}B, \quad M_K^2 = (\hat{m} + m_s)B, \quad M_{\eta}^2 = \frac{2}{3}(\hat{m} + 2m_s)B.$$
(15)

Meson masses obviously satisfy the Gell-Mann–Oakes– Renner and the Gell-Mann–Okubo relation $3M_{\eta}^2 + M_{\pi}^2 = 4M_K^2$. In the evaluation we use the following set of QCD parameters [24]: $\hat{m} = 7$ MeV, $m_s/\hat{m} = 25$, $B = M_{\pi^+}^2/(2\hat{m}) = 1.4$ GeV. For the pion decay constant *F*, defined in the chiral limit, we take the value of 88 MeV, calculated in ChPT [5].

Our next approximation consists of a linearizing the resulting Lagrangian \mathcal{L}_{full} with respect to the chiral field $\hat{\Phi}$:

$$U = \exp\left[i\frac{\hat{\Phi}}{F}\right] \simeq 1 + i\frac{\hat{\Phi}}{F} + o\left(\frac{\hat{\Phi}}{F}\right).$$
(16)

In other words, we treat Goldstone fields as small fluctuations around the three-quark (3q) core. Considering the 3q baryon core as a zeroth-order approximation we take into account mesonic degrees of freedom by a *perturbative expansion*, i.e., expansion in terms of the coupling 1/*F*. In the evaluation of matrix elements we restrict ourselves to the one-loop approximation, i.e., we perform our calculations to order of accuracy $o(1/F^2)$. Finally, with the use of ansatz (16) our full Lagrangian \mathcal{L} [see Eqs. (13) and (14)] reduces to the effective linearized Lagrangian

$$\mathcal{L}_{eff}(x) = \bar{\psi}(x) [i\partial - S(r) - \gamma^0 V(r)] \psi(x) + \frac{1}{2} (\partial_\mu \Phi_i)^2 - \bar{\psi}(x) S(r) i \gamma^5 \frac{\hat{\Phi}}{F} \psi(x) + \mathcal{L}_{\chi SB}(x).$$
(17)

To describe the properties of baryons as bound states of quarks surrounded by a meson cloud we have to formulate perturbation theory. First, let us specify how we calculate the baryon mass spectrum (here and in the following we restrict considerations only to the nucleon and Δ isobar). In our approach the masses (energies) of the three-quark core for the nucleon and Δ isobar are degenerate and are related to the single quark energy \mathcal{E}_0 by $E_0 = m_N = m_\Delta = 3 \cdot \mathcal{E}_0$ (In this paper we do not discuss the removal of the spurious contribution to the baryon mass arising from the center-of-mass motion of the bound state). The energy of the bound state E_0 satisfies the eigenequation $H_0 |\phi_0\rangle = E_0 |\phi_0\rangle$ with H_0 being the free quark Hamiltonian and $|\phi_0
angle$ being the unperturbed three-quark state. The single quark ground state energy \mathcal{E}_0 is obtained from the Dirac equation for the ground state quark wave function (w.f.) $u_0(x)$ with

$$\left[-i\vec{\alpha}\vec{\nabla}+\beta S(r)+V(r)-\mathcal{E}_{0}\right]u_{0}(\vec{x})=0,$$
(18)

where we use the standard spinor algebra notation: $\vec{\alpha} = \gamma^0 \vec{\gamma}$ and $\beta = \gamma^0$. The quark w.f. $u_0(\vec{x})$ belongs to the basis of potential eigenstates (including excited quark and antiquark solutions) used for expansion of the quark field operator $\psi(x)$. Here we restrict the expansion to the ground state contribution with $\psi(x) = b_0 u_0(x) \exp(-i\mathcal{E}_0 t)$, where b_0 is the corresponding single quark annihilation operator. In Eq. (18) we drop the current quark mass to simplify our calculational technique. Instead we consider the quark mass term as a small perturbation [24]. Inclusion of a finite current quark mass leads to a displacement of the single quark energy which, by Eq. (5), is relevant for the calculation of the sigma term; a quantity which vanishes in the chiral limit. On the other hand, the effect of a finite current mass on observables which survive in the chiral limit, like magnetic moments, charge radii, etc., are quite negligible.

In general, for a given form of the potentials S(r) and V(r) the Dirac Eq. (18) can be solved numerically. Here, for the sake of simplicity, we use a variational *Gaussian ansatz* [25] for the quark wave function given by the analytical form

$$u_0(\vec{x}) = N \exp\left[-\frac{\vec{x}^2}{2R^2}\right] \begin{pmatrix} 1\\ i\rho\vec{\sigma}\vec{x}/R \end{pmatrix} \chi_s \chi_f \chi_c, \quad (19)$$

where $N = [\pi^{3/2}R^3(1+3\rho^2/2)]^{-1/2}$ is a constant fixed by the normalization condition $\int d^3x u_0^{\dagger}(x) u_0(x) \equiv 1$; χ_s , χ_f , χ_c are the spin, flavor and color quark wave functions, respectively. Our Gaussian ansatz contains two model parameters: the di-

mensional parameter *R* and the dimensionless parameter ρ . The parameter ρ can be related to the axial coupling constant g_A calculated in zeroth-order (or 3q-core) approximation:

$$g_A = \frac{5}{3} \left(1 - \frac{2\rho^2}{1 + \frac{3}{2}\rho^2} \right).$$
(20)

For $\rho = 0$ we obtain the nonrelativistic value of $g_A = 5/3$. The parameter *R* can be physically understood as the mean radius of the three-quark core and is related to the charge radius $\langle r_E^2 \rangle^P$ of the proton in the zeroth-order approximation as

$$\langle r_E^2 \rangle_{3q-core}^P = \int d^3x u_0^{\dagger}(\vec{x}) \vec{x^2} u_0(\vec{x}) = \frac{3R^2}{2} \frac{1 + \frac{5}{2}\rho^2}{1 + \frac{3}{2}\rho^2}.$$
(21)

In our calculations we use the value $g_A = 1.25$ as obtained in ChPT [5]. Therefore we have only one free parameter *R*. In the numerical evaluation *R* is varied in the region from 0.55 to 0.65 fm corresponding to a change of $\langle r_E^2 \rangle_{3q-core}^P$ in the region from 0.5 to 0.7 fm². The use of the Gaussian ansatz (19) in its exact form restricts the scalar confinement potential *S*(*r*) to

$$S(r) = \frac{1 - 3\rho^2}{\underbrace{2\rho R}_{=M}} + \underbrace{\frac{\rho}{2R^3}}_{=c} r^2 = M + cr^2, \qquad (22)$$

expressed in terms of the parameters R and ρ . The constant part of the scalar potential M can be interpreted as the constituent mass of the quark, which is simply a displacement of the current quark mass due to the potential S(r). The parameter c is the coupling defining the radial (quadratic) dependence of the scalar potential. Numerically, for our set of parameters, we get $M = 230 \pm 20$ MeV and $c = 0.08 \pm 0.01$ GeV³. These and the following error bars in our results correspond to the variation of the parameter R. For the vector potential we get the following expression:

$$V(r) = \mathcal{E}_0 - \frac{1+3\rho^2}{2\rho R} + \frac{\rho}{2R^3}r^2,$$
 (23)

where the single quark energy \mathcal{E}_0 is a free parameter in the Gaussian ansatz. For our purposes (calculation of the sigma terms) the magnitude of the quark energy in zeroth-order \mathcal{E}_0 does not influence the result. Nevertheless, in view of Eq. (5), the shift of \mathcal{E}_0 due to a finite quark mass *m* will contribute to the meson-baryon sigma terms. The displacement of the quark energy in an expansion to first order in *m*, with the use our Gaussian ansatz, is

$$\mathcal{E}_0 \to \mathcal{E}_0(m) = \mathcal{E}_0 + \gamma m + o(m), \qquad (24)$$

where γ is the relativistic reduction factor

$$\gamma = \frac{1 - \frac{3}{2}\rho^2}{1 + \frac{3}{2}\rho^2} = \frac{9}{10}g_A - \frac{1}{2}.$$
 (25)

In the nonrelativistic limit ($\rho = 0$) we get $\gamma = 1$ and therefore the corresponding displacement is equal to the current quark mass. In the relativistic picture, for $\rho^2 > 0$, we definitely have $\gamma < 1$. Particularly for our choice of $g_A = 1.25$ we get $\gamma = 5/8$.

In our approach the PCAC (partial conservation of axialvector current) requirement and consequently the Goldberger-Treiman relation are fulfilled. In other words, the expectation value of the pseudoscalar isovector density J_i = $(1/F)S(r)\bar{\psi}i\gamma_5\tau_i\psi$ between unperturbed 3q states $|\phi_0\rangle$ states leads to the pion-nucleon constant $G_{\pi NN}$ and we arrive at the Goldberger-Treiman relation between the couplings $G_{\pi NN}$ and g_A with

$$G_{\pi NN} = \frac{5m_N}{3F} \left(1 - \frac{2\rho^2}{1 + \frac{3}{2}\rho^2} \right) \equiv \frac{m_N}{F} g_A.$$
(26)

The same condition holds even for their form factors. The analytical expression for the pion-nucleon form factor in the chiral limit is given by

$$G_{\pi NN}(Q^2) = \frac{m_N}{F} g_A(Q^2) = \frac{m_N}{F} g_A F_{\pi NN}(Q^2), \quad (27)$$

where Q^2 is the squared Euclidean momentum of the pion and $F_{\pi NN}(Q^2)$ is the πNN form factor normalized to unity at zero recoil $Q^2 = 0$:

$$F_{\pi NN}(Q^2) = \exp\left(-\frac{Q^2 R^2}{4}\right) \left\{ 1 + \frac{Q^2 R^2}{8} \left(1 - \frac{5}{3g_A}\right) \right\}.$$
(28)

Following the Gell-Mann and Low theorem [26] the energy shift ΔE_0 of the three-quark ground state due to the interaction with Goldstone mesons is given by the expression

$$\Delta E_0 = \langle \phi_0 | \sum_{i=1}^{\infty} \frac{(-i)^n}{n!} \\ \times \int i \,\delta(t_1) d^4 x_1 \cdots d^4 x_n T[\mathcal{H}_I(x_1) \cdots \mathcal{H}_I(x_n)] | \phi_0 \rangle_c \,,$$
(29)

where

$$\mathcal{H}_{I}(x) = \bar{\psi}(x)i\gamma^{5}\frac{\hat{\Phi}(x)}{F}S(r)\psi(x)$$
(30)

is the interaction Hamiltonian and subscript "*c*" refers to contributions from connected graphs only. We evaluate Eq. (29) at one loop with $o(1/F^2)$ using Wick's theorem and the appropriate propagators. For the quark field we use a Feyn-

man propagator for a fermion in a binding potential. By restricting the summation over intermediate quark states to the ground state we get

$$iG_{\psi}(x,y) = \langle \phi_0 | T\{\psi(x)\overline{\psi}(y)\} | \phi_0 \rangle \rightarrow u_0(\overline{x})\overline{u}_0(\overline{y})$$
$$\times \exp[-i\mathcal{E}_0(x_0 - y_0)] \theta(x_0 - y_0). \tag{31}$$

For meson fields we use the free Feynman propagator for a boson field with

$$i\Delta_{ij}(x-y) = \langle 0|T\{\Phi_i(x)\Phi_j(y)\}|0\rangle$$
$$= \delta_{ij} \int \frac{d^4k}{(2\pi)^4 i} \frac{\exp[-ik(x-y)]}{M_{\Phi}^2 - k^2 - i\epsilon}.$$
 (32)

III. MESON-BARYON SIGMA TERMS IN THE PCQM

The scalar density operators S_i^{PCQM} (i=u,d,s), relevant for the calculation of the meson-baryon sigma terms in the PCQM, are defined as the partial derivatives of the model χ SB Hamiltonian $\mathcal{H}_{\chi SB} = -\mathcal{L}_{\chi SB}$ with respect to the current quark mass of *i*th flavor m_i . Note that the nondiagonal term in $\mathcal{H}_{\chi SB}$ which is proportional to $(m_u - m_d)\pi^0\eta$ vanishes because we apply the isospin limit with $m_u = m_d$. Here we obtain

$$S_i^{PCQM} \doteq \frac{\partial \mathcal{H}_{\chi SB}}{\partial m_i} = S_i^{val} + S_i^{sea}, \qquad (33)$$

where S_i^{val} is the set of valence-quark operators coinciding with the ones obtained from the QCD Hamiltonian (2)

$$S_u^{val} = \overline{u}u, \quad S_d^{val} = \overline{d}d, \quad S_s^{val} = \overline{s}s.$$
 (34)

The set of sea-quark operators S_i^{sea} arises from the pseudoscalar meson mass term [due to Eq. (15)] with

$$S_{u}^{sea} = B\left\{\pi^{+}\pi^{-} + \frac{\pi^{0}\pi^{0}}{2} + K^{+}K^{-} + \frac{\eta^{2}}{6}\right\}, \quad (35)$$
$$S_{d}^{sea} = B\left\{\pi^{+}\pi^{-} + \frac{\pi^{0}\pi^{0}}{2} + K^{0}\bar{K}^{0} + \frac{\eta^{2}}{6}\right\},$$
$$S_{s}^{sea} = B\left\{K^{+}K^{-} + K^{0}\bar{K}^{0} + \frac{2}{3}\eta^{2}\right\}.$$

To calculate meson-baryon sigma terms, equivalent to the definition of Eq. (4), we perform the perturbative expansion for the matrix element of the scalar density operator S_i^{PCQM} between unperturbed 3q-core states:

$$\langle \phi_0 | \sum_{n=0}^{\infty} \frac{(-i)^n}{n!} \int d^4 x_1 \cdots \int d^4 x_n T[\mathcal{H}_I(x_1) \cdots \mathcal{H}_I(x_n) \\ \times S_i^{PCQM}] | \phi_0 \rangle_c .$$
 (36)

In the evaluation of Eqs. (29) and (36) we project the matrix elements on the respective baryon states. Baryon wave functions for nucleon and delta states are conventionally set up



FIG. 1. Diagrams contributing to the baryon energy shift: meson cloud (a) and meson exchange diagram (b).

by the product of the SU(6) spin-flavor w.f. (wave function) and SU(3)_c color w.f. (see details in Ref. [27]), where the nonrelativistic single quark spin w.f. is replaced by the relativistic ground state solution of Eq. (19).

The diagrams that contribute to the energy shift ΔE_0 at one loop are shown in Fig. 1: meson cloud [Fig. 1(a)] and meson exchange diagram [Fig. 1(b)]. The explicit expressions for the nucleon and Δ -isobar masses including oneloop corrections are given by

$$m_{B} = 3(\mathcal{E}_{0} + \gamma \hat{m}) + d_{B}^{\pi} \Pi(M_{\pi}^{2}) + d_{B}^{K} \Pi(M_{K}^{2}) + d_{B}^{\eta} \Pi(M_{\eta}^{2})$$
(37)

with
$$d_N^{\pi} = \frac{171}{400}$$
, $d_{\Delta}^{\pi} = \frac{11}{19} d_N^{\pi}$, $d_B^{K} = \frac{6}{19} d_N^{\pi}$,
 $d_N^{\eta} = \frac{1}{57} d_N^{\pi}$, $d_{\Delta}^{\eta} = \frac{5}{57} d_N^{\pi}$,

where d_B^{Φ} with $\Phi = \pi, K$ or η and B = N or Δ are the recoupling coefficients defining the partial contribution of the π , K, and η -meson cloud to the energy shift of the nucleon and Δ isobar, respectively. The self-energy operators $\Pi(M_{\Phi}^2)$, corresponding to meson cloud contributions with definite flavor, differ only in their value for the meson mass and are given by

$$\Pi(M_{\Phi}^2) = -\left(\frac{g_A}{\pi F}\right)^2 \int_0^\infty \frac{dpp^4}{p^2 + M_{\Phi}^2} F_{\pi NN}^2(p^2), \qquad (38)$$

where $p = |\vec{p}|$ is the absolute value of the three momentum of the meson. Using Eqs. (37) and (38) we obtain the expressions for the baryon (nucleon and Δ -isobar) masses

$$m_{B} = \underbrace{3\mathcal{E}_{0} + \Pi(0) \cdot \sum_{\Phi = \pi, K, \eta} d_{B}^{\Phi}}_{= m_{B}}$$

$$+ \underbrace{3\gamma\hat{m} + \sum_{\Phi = \pi, K, \eta} d_{B}^{\Phi} \cdot [\Pi(M_{\Phi}^{2}) - \Pi(0)]}_{= \Delta m_{B}}$$

$$= \stackrel{0}{m_B} + \Delta m_B \,. \tag{39}$$



FIG. 2. Diagrams contributing to the meson-baryon sigma terms: tree diagram (a), meson cloud diagram (b), and meson exchange diagram (c). Insertion of the scalar density operator is depicted by the symbol \bigvee .

The baryon mass in the chiral limit $(\hat{m}, m_s \rightarrow 0)$ is m_B^0 . Meson loops also contribute to m_B^0 , since in the chiral limit quarks interact with massless mesons. In the calculations of the one-loop meson contributions to the baryon masses we do not take into account the modification of the quark w.f. due to a small but finite current quark mass. This effect is strongly suppressed and therefore, as mentioned above, we work at the order of accuracy $o(\hat{m}, m_s, 1/F^2)$.

At the same level of accuracy $o(\hat{m}, m_s, 1/F^2)$ we calculate the meson-baryon sigma terms using Eq. (36). The following diagrams contribute to the sigma terms up to the one-loop level: tree level diagram [Fig. 2(a)] with the insertion of the valence-quark scalar density S_i^{val} into the quark line, meson cloud [Fig. 2(b)], and meson exchange diagrams [Fig. 2(c)] with insertion of the sea-quark scalar density S_i^{sea} to the meson line. We neglect diagrams with an insertion of S_i^{val} to the quark lines in the second-order graphs [like in Figs. 1(a) and (b)]. These terms are proportional to \hat{m}/F^2 and therefore are of higher order. First, we consider the πN sigma term. The analytical expression for this quantity is derived as

where the first term of the right-hand side of Eq. (40) corresponds to the valence quark, the second to the sea quark contribution. The vertex function $\Gamma(M_{\Phi}^2)$, as derived by use of Eq. (36), can be related to the partial derivative of the self-energy operator $\Pi(M_{\Phi}^2)$ with respect to the nonstrange current quark mass \hat{m} with

(

$$\Gamma(M_{\Phi}^2) = \hat{m} \frac{\partial}{\partial \hat{m}} \Pi(M_{\Phi}^2).$$
(41)

The derivative $\hat{m}\partial/\partial \hat{m}$ is equivalent to the one with respect to the meson masses [1]:

$$\hat{m}\frac{\partial}{\partial \hat{m}}\Pi(M_{\Phi}^2) = M_{\pi}^2 \left(\frac{\partial}{\partial M_{\pi}^2} + \frac{1}{2}\frac{\partial}{\partial M_K^2} + \frac{1}{3}\frac{\partial}{\partial M_{\eta}^2}\right)\Pi(M_{\Phi}^2).$$
(42)

Using Eqs. (39)–(41) we directly prove the Feynman-Hellmann theorem of Eq. (5); hence our approach is consistent to order of accuracy $o(\hat{m}, m_s, 1/F^2)$. Our next point is the analysis of the chiral expansion for the πN sigma term $\sigma_{\pi N}^{\pi}$ when we restrict to the two-flavor picture, i.e., we take into account only pion contributions. The expansion for $\sigma_{\pi N}^{\pi}$ in the PCQM is given by an analytic expression in terms of fundamental constants (the couplings g_A and F, the quark and pion masses \hat{m} and M_{π})

$$\sigma_{\pi N}^{\pi} = \frac{3\gamma}{2B} M_{\pi}^{2} + \frac{d_{N}^{\pi}}{(2\pi)^{3/2}} \left(\frac{g_{A}}{F}\right)^{2} \frac{M_{\pi}^{2}}{R}$$

$$\times \sum_{N=0}^{\infty} \frac{(-M_{\pi}R)^{N}}{2^{N/2}} \frac{\Gamma\left(\frac{1}{2}\right)}{\Gamma\left(\frac{N+1}{2}\right)} (N+2)\mathcal{F}_{N}(\gamma),$$
(43)

where

$$\mathcal{F}_{N}(\gamma) = 1 + \left(\frac{1-\gamma}{1+2\gamma}\right) \frac{N-1}{2} + \left(\frac{1-\gamma}{1+2\gamma}\right)^{2} \frac{(N-1)(N-3)}{16}$$

is a polynomial in the relativistic factor γ introduced before. Explicitly, up to order $o(M_{\pi}^5)$ we obtain

$$\sigma_{\pi N}^{\pi} = k_2 M_{\pi}^2 + \frac{3}{2} k_3 M_{\pi}^3 + 2k_4 M_{\pi}^4 + \frac{5}{2} k_5 M_{\pi}^5 + o(M_{\pi}^5),$$
(44)

where k_i are the expansion coefficients evaluated as

$$k_{2} = \frac{3\gamma}{2B} + \frac{171}{200} \left(\frac{g_{A}}{F}\right)^{2} \frac{1+2\xi-3\xi^{2}}{(2\pi)^{3/2}R},$$

$$k_{3} = -\frac{171}{800\pi} \left(\frac{g_{A}}{F}\right)^{2}, \quad k_{4} = \frac{171}{200} \left(\frac{g_{A}}{F}\right)^{2} R \frac{1-2\xi-\xi^{2}}{(2\pi)^{3/2}},$$

$$k_{5} = -\frac{171}{1600} \left(\frac{g_{A}}{F}\right)^{2} \frac{R^{2}}{\pi} (1-4\xi),$$

$$\xi = \frac{1-\gamma}{4(1+2\gamma)}.$$
(45)

In the present approach the coefficient k_3 defines the leading nonanalytic contribution (LNAC) to $\sigma_{\pi N}$. Both, nucleon and Δ -isobar degrees of freedom as intermediate states in the meson-loop diagrams contribute to the coefficient k_3 . The nucleonic contribution to the coefficient k_3 is $k_3^N =$ $-3g_A^2/(32\pi F^2)$. The Δ isobar also contributes to k_3 with $k_3^{\Delta} = 32/25k_3^N$. In total, our LNA coefficient k_3 results in k_3 $=k_3^N+k_3^{\Delta}=-171g_A^2/(800\pi F^2)$ as also obtained in a similar quark model approach [28]. One should remark that our LNAC disagrees with a similar quantity given in effective hadronic approaches [7,24,29,30] and in the cloudy bag model [31]. In later approaches the unperturbed nucleon and Δ states are not degenerate. Therefore only the loop diagram with πN as an intermediate state contributes to the LNAC which is equal to $-3g_A^2/(32\pi F^2)$ and coincides with our k_3^N . The $\pi \Delta$ loop diagram contributes to the next-to-leading order nonanalytic term.

Numerically, the contribution of the valence quarks to the πN sigma term is 13.1 MeV, the contribution of sea quarks at order $o(M_{\pi}^2)$ is 66.9 ± 5.7 MeV. Higher order contributions of the sea quarks are $(3k_3/2)M_{\pi}^3 = -57.2$ MeV, $2k_4M_{\pi}^4 = 27.7\pm2.3$ MeV, $(5k_5/2)M_{\pi}^5 = -10\pm1.7$ MeV, and $o(M_{\pi}^5) \approx 2.8$ MeV. Therefore, at order $o(M_{\pi}^5)$ in the two-flavor picture we have the following result for the πN sigma term $\sigma_{\pi N}^{\pi}$ [again, superscript π refers to the SU(2) flavor picture] of

$$\sigma_{\pi N}^{\pi(5)} = 40.5 \pm 5 \text{ MeV},$$
 (46)

compared to the full calculation of the πN sigma term of Eq. (43) of

$$\sigma_{\pi N}^{\pi} = 43.3 \pm 4.4 \text{ MeV.}$$
 (47)

As we mentioned before, the error bars are due to a variation of the range parameter R of the quark wave function (19) from 0.55 to 0.65 fm.

To complete our estimate of the πN sigma term we evaluate the additional contributions of kaon and η -meson loops, $\sigma_{\pi N}^{K}$ and $\sigma_{\pi N}^{\eta}$, where superscripts *K* and η refer to the respective meson cloud contribution. These terms are significantly suppressed relative to the pion cloud and to the valence quark contributions. For illustration, we give the expressions for r_K (ratio of kaon to pion cloud contributions) and r_{η} (ratio of η meson to pion cloud contributions):

$$r_{K} = \frac{\sigma_{\pi N}^{K}}{\sigma_{\pi N}^{\pi} - 3\,\gamma \hat{m}} = \frac{d_{N}^{K}}{2d_{N}^{\pi}} \cdot \frac{\int_{0}^{\infty} \frac{dpp^{4}}{(p^{2} + M_{K}^{2})^{2}} F_{\pi N N}^{2}(p^{2})}{\int_{0}^{\infty} \frac{dpp^{4}}{(p^{2} + M_{\pi}^{2})^{2}} F_{\pi N N}^{2}(p^{2})}, \qquad (48)$$

$$r_{\eta} = \frac{\sigma_{\pi N}^{\eta}}{\sigma_{\pi N}^{\pi} - 3\,\gamma \hat{m}} = \frac{d_{N}^{\eta}}{3d_{N}^{\pi}} \cdot \frac{\int_{0}^{\infty} \frac{dpp^{4}}{(p^{2} + M_{\eta}^{2})^{2}} F_{\pi N N}^{2}(p^{2})}{\int_{0}^{\infty} \frac{dpp^{4}}{(p^{2} + M_{\eta}^{2})^{2}} F_{\pi N N}^{2}(p^{2})}.$$
 (49)

The energy denominators in Eqs. (48) and (49) ensure that the contributions of the kaon and η -meson cloud to the πN sigma term are negligible when compared to the pionic one. The same conclusion regarding the suppression of *K* and η -meson loops was obtained in the cloudy bag model [32]. Numerically, kaon and η -meson cloud contributions are $\sigma_{\pi N}^{K} = 1.7 \pm 0.4$ MeV and $\sigma_{\pi N}^{\eta} = 0.023 \pm 0.006$ MeV with $r_{K} = 5.6 \pm 1.6\%$ and $r_{\eta} = 0.08 \pm 0.02\%$.

For the πN sigma term we have the following final value:

$$\sigma_{\pi N} = \sum_{\Phi = \pi, K, \eta} \sigma_{\pi N}^{\Phi} = 45 \pm 5 \quad \text{MeV.}$$
(50)

Our result for the πN sigma term is in perfect agreement with the value of $\sigma_{\pi N} \approx 45$ MeV deduced by Gasser, Leutwyler, and Sainio [33] using dispersion-relation techniques and exploiting the chiral symmetry constraints.

In our opinion, a meaningful description of the πN sigma term should be based on the following guide lines: chiral symmetry constraints, fulfillment of low-energy theorems, consistency with the Feynman-Hellmann theorem, and proper treatment of sea quarks, that is meson cloud contributions. The model developed here fulfils these aspects. In the literature there has been considerable effort to determine the πN sigma term in the framework of different approaches: effective field theories, lattice QCD, QCD sum rules, quark models, soliton-type approaches, etc. In HBChPT [12–14] the πN sigma term is used as an input parameter to fix the couplings in the effective Lagrangian. Then the KN sigma terms are predicted (see discussion later on). A detailed analysis of the πN sigma term was done in the cloudy bag model (CBM) [32,34], which is similar to our approach. Numerical results obtained in the CBM [32] and the PCQM are quite close. For example, in both calculations the dominant contribution to the πN sigma term is due to the pion cloud: 22.7 MeV (in CBM) and ≈ 30 MeV (in PCQM). Contributions of kaon and η -meson loops are significantly suppressed in both approaches. A determination of $\sigma_{\pi N}$ in lattice QCD has been undertaken by several groups using different techniques: direct calculation of the scalar density matrix element of the nucleon and use of the Feynman-Hellmann theorem (see detailed discussion in Ref. [35]). As correctly pointed out in Ref. [35], the main disadvantage of all lattice approaches based on the calculation of the scalar matrix element of the nucleon is that latter quantity and the current quark mass \hat{m} are not renormalization group (RG) invariant. This leads to uncertainties in the evaluation of the πN sigma term which should be RG invariant. In a recent paper [35] (see also Ref. [36]) the πN sigma term was calculated by means of the Feynman-Hellmann theorem and using present SU(2) lattice QCD data for the nucleon mass as a function of the current quark mass. The main idea of Ref. [35] is to include the meson cloud contributions properly as based on chiral symmetry constraints; the corresponding expression for the nucleon mass as function of \hat{m} includes meson loop contribution (πN and $\pi \Delta$) with correct leading and next-toleading order nonanalytic behavior [36]. Applying this extrapolation function to the SU(2) lattice data results in a value of $\sigma_{\pi N} = 45 \div 55$ MeV. This result is quite close to our estimate of the πN sigma term in the two-flavor case [see Eq. (47)]. Another example for the dominance of the pion cloud contributions to the πN sigma term are the solitontype quark models: the chiral quark soliton model [37] and the confining chiral soliton model [38]. Their results $\sigma_{\pi N}$ =54.3 MeV [37] and $\sigma_{\pi N}$ =30–40 MeV [38] are also in qualitative agreement with the previously stated approaches (cloudy bag model [32], lattice QCD [35]) and the result obtained here.

One of the advantages of our approach is the possibility to estimate free coupling constants in the effective Lagrangians of ChPT, Baryon ChPT (BChPT), and HBChPT. For example, in BChPT [7] and HBChPT [6] the quadratic term $(-4c_1M_{\pi}^2)$ in the chiral expansion of the πN sigma term contains the unknown coupling c_1 of the effective ChPT Lagrangian [7]. Using Eq. (45) we predict $c_1 = -k_2/4 =$ -1.16 ± 0.1 GeV⁻¹. This prediction is in good agreement with the value of $c_1 = -0.925$ GeV⁻¹ determined by Becher and Leutwyler [39] from a fit of the elastic πN scattering amplitude at threshold to data of the Karlsruhe partial wave analysis (KA86 data).

Next we discuss our prediction for the strangeness content of the nucleon y_N which is defined in the PCQM as

$$y_N = \frac{2\langle p|S_s^{PCQM}|p\rangle}{\langle p|S_u^{PCQM} + S_d^{PCQM}|p\rangle}.$$
(51)

The direct calculation of the strange-quark scalar density $\langle p|S_s^{PCQM}|p\rangle$ using Eq. (36) is completely consistent with the indirect one applying the Feynman-Hellmann theorem $\langle p|S_s^{PCQM}|p\rangle = \partial m_N / \partial m_s$ with

$$\langle p | S_s^{PCQM} | p \rangle = \frac{\partial m_N}{\partial m_s} = B \left[\frac{\partial}{\partial M_K^2} + \frac{4}{3} \frac{\partial}{\partial M_\eta^2} \right] m_N$$

$$= \left(\frac{g_A}{\pi F} \right)^2 B \int_0^\infty dp p^4 F_{\pi NN}^2 (p^2)$$

$$\times \left[\frac{d_N^K}{(p^2 + M_K^2)^2} + \frac{4}{3} \frac{d_N^\eta}{(p^2 + M_\eta^2)^2} \right].$$
(52)

Combining Eqs. (40) and (52) we derive for y_N following expression:

$$y_{N} = \frac{\left(\frac{g_{A}}{\pi F}\right)^{2} B \int_{0}^{\infty} dp p^{4} F_{\pi NN}^{2}(p^{2}) \left[\frac{d_{K}}{(p^{2} + M_{K}^{2})^{2}} + \frac{4}{3} \frac{d_{\eta}}{(p^{2} + M_{\eta}^{2})^{2}}\right]}{\frac{3}{2} \gamma + \left(\frac{g_{A}}{\pi F}\right)^{2} B \int_{0}^{\infty} dp p^{4} F_{\pi NN}^{2}(p^{2}) \left[\frac{d_{\pi}}{(p^{2} + M_{\pi}^{2})^{2}} + \frac{1}{2} \frac{d_{K}}{(p^{2} + M_{K}^{2})^{2}} + \frac{1}{3} \frac{d_{\eta}}{(p^{2} + M_{\eta}^{2})^{2}}\right]},$$
(53)

with the numerical value of

(

a

$$y_N = 0.076 \pm 0.012.$$
 (54)

The small value of y_N in our model is due to the suppressed contributions of kaon and η -meson clouds. Our prediction for y_N is smaller than the value $y_N \approx 0.2$ obtained in Ref. [33] from an analysis of experimental data on πN phase shifts. On the other hand, our prediction is quite close to the result obtained in the Skyrme model $y_N \approx 0.058$ [40] and in the cloudy bag model $y_N \approx 0.05$ [32]. A revisited prediction of HBChPT for y_N gives $y_N = 0.25 \pm 0.05$ (without inclusion of the decuplet) and $y_N = 0.20 \pm 0.12$ (taking into account the strange content of the nucleon in lattice approaches gives larger values for $y_N = 0.36 \pm 0.03$ [41] and $y_N = 0.59 \pm 0.13$ [42], which imply a big contribution of the strange quark sea to the nucleon mass.

Next, we present our results for the KN sigma terms:

$$\sigma_{KN}^{u} \equiv \sigma_{KN}^{(1)} = 340 \pm 37 \text{ MeV}, \quad \sigma_{KN}^{d} = 284 \pm 37 \text{ MeV},$$
(55)

$$\sigma_{KN}^{(2)} = 228 \pm 37 \text{ MeV}, \quad \sigma_{KN}^{I=1} = 28 \text{ MeV},$$

and $\sigma_{KN} \equiv \frac{\sigma_{KN}^{u} + \sigma_{KN}^{d}}{2} = 312 \pm 37 \text{ MeV}.$

Comparative results for the *KN* sigma term are less abundant than for the πN sigma terms. We cite recent (revisited) results of HBChPT (taking into account the decuplet contribution) with $\sigma_{KN}^{(1)}=380\pm40$ MeV, $\sigma_{KN}^{(2)}=250\pm30$ MeV [14], results of lattice QCD with $\sigma_{KN}=(\sigma_{KN}^u+\sigma_{KN}^d)/2=362$ ± 13 MeV [41], and predictions of the Nambu-Jona-Lasinio model with $\sigma_{KN}=425$ MeV (with an error bar of 10–15%) [43]. Hopefully, future DA Φ NE experiments at Frascati [44] will allow for a determination of the *KN* sigma terms and hence for a better knowledge of the strangeness content of the nucleon.

Finally, for the sake of completeness we present the set of our predictions for the Δ -isobar condensates and ηN sigma term:

$$\sigma_{\pi\Delta} = 32 \pm 3$$
 MeV, $y_{\Delta} = 0.12 \pm 0.02$,
 $\sigma_{\eta N} = 72 \pm 16$ MeV, (56)
 $\sigma_{K\Delta}^{u} = 262 \pm 26$ MeV,
 $\sigma_{K\Delta}^{d} = 206 \pm 26$ MeV, $\sigma_{\eta \Delta} = 75 \pm 17$ MeV.

IV. SUMMARY AND CONCLUSIONS

In conclusion, we have evaluated the meson-baryon sigma terms using a perturbative chiral quark model based on an

effective chiral Lagrangian. The Lagrangian describes baryons as bound states of three valence quarks surrounded by a cloud of pseudoscalar mesons as dictated by the chiral symmetry requirement. The calculated quantities contain only one model parameter R, which is related to the radius of the three-quark core, and are otherwise expressed in terms of fundamental parameters of low-energy hadron physics: axial coupling constant g_A , weak pion decay constant F, normalized strong pion-nucleon form factor $F_{\pi NN}$, and set of QCD parameters (current quark masses \hat{m} and m_s and quark condensate parameter B). Predictions are given for a variation of the free parameter R in a quite wide physical region from 0.55 to 0.65 fm corresponding to $\langle r_E^2 \rangle_{3q-core}^P$ ranging from 0.5 to 0.7 fm². Our result for the πN sigma term is in perfect agreement with the value obtained by Gasser, Leutwyler, and Sainio [33] and with the results of other theoretical approaches (cloudy bag model [32], chiral quark soliton model [37], lattice QCD [35], etc.) where the pion cloud

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contribution properly was taken into account. Analyses of the strangeness content of the nucleon, kaon-baryon, and etabaryon sigma terms were done in detail. We compare our predictions to results of other theoretical approaches. Ongoing experimental efforts aim to allow for a reliable extraction of these quantities, and hence to fix the scalar strange-quark density in the nucleon. To solidify the model approach we mention that with the same values of the free model parameter R we recently obtained [16] a quite reasonable description of static properties and electromagnetic form factors of the nucleon.

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