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Chiral symmetry and the Higgs-boson–nucleon coupling

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The chiral-symmetry-breaking σ term as extracted from the pion-nucleon phase-shift and dispersion analysis implies that the Higgs boson coupling to the nucleon is dominated by the strange quarks. For $\sigma_{\pi N} \simeq 60$ MeV the strange-quark contribution is an order of magnitude larger than that of any other flavor. This significantly increases the Higgs-boson–nucleon coupling from its chiral-symmetric value. Some of the phenomenological consequences are briefly discussed. In particular I reanalyze the low-energy nuclear and atomic experimental evidence against the existence of a light scalar boson, and compare the bounds to those originally derived with the expectation that the nucleon consists mainly of the up and down quarks. A summary review of the various issues concerning a large πN σ term is also provided.

I. INTRODUCTION

The pion-nucleon σ terms can be expressed as the quark scalar densities taken between the nucleon state (at zero momentum transfer):

$$\sigma_{\pi N} = \hat{m} \langle N \mid (\bar{u}u + \bar{d}d) \mid N \rangle \ , \tag{1}$$

with $\hat{m} = \frac{1}{2}(m_u + m_d)$. Current algebra[1] relates this quantity directly to the pion-nucleon amplitude in the off-shell soft-pion limit. In addition, in the lowest order of chiral perturbation theory, it is related to the isospin-even on-shell amplitude at an unphysical energy and momentum-transfer point[2] of $s = M_N^2$ and $t = 2m_\pi^2$:

$$T_{\pi N}^{(+)}(M_N^2, 2m_\pi^2) \simeq \sigma_{\pi N} / f_\pi^2 \ , \tag{2}$$

where f_π is the pion decay constant. This πN amplitude can be extracted from phase-shift analysis and dispersion calculations. The double extrapolation, although very delicate, has been performed by many groups.[3] In addition, ever since the extremely accurate measurements of the low-energy pion-nucleon scattering cross sections by Carter and co-workers,[4] all these extractions have yielded consistent results, all in the range between 50 and 70 MeV (Ref. 5). The most systematic and complete analysis has been by the Karlsruhe group. A result of $f_\pi^2 T_{\pi N}^{(+)}(M_N^2, 2m_\pi^2) \simeq 64$ MeV has been deduced.[6] The higher-order corrections to Eq. (2) are of the order m_π^3 and $m_\pi^4 \ln m_\pi^2$ and they reduce the σ term by no more than 5 MeV (Ref. 7). Thus for the analysis in this paper we quote an "experimental value" of

$$\sigma_{\pi N} \simeq 60 \text{ MeV} \ . \tag{3}$$

What is the theoretical implication of this result? The initial motivations were to test QCD as a concrete realization of the Gell-Mann–Oakes–Renner chiral-symmetry-breaking model in which the SU(3)×SU(3)-breaking term transforms according to the $(3,3^*)+(3^*,3)$ representation.[8] Nevertheless, $\sigma_{\pi N}$ cannot be calculated based on symmetry considerations alone; dynamical assumptions, in particular, with respect to the strange-quark content of the nucleon, must be invoked.[9] Let us review this analysis. Ignoring the tiny isospin-violating term of $(m_u - m_d)(\bar{u}u - \bar{d}d)$, the chiral-symmetry-breaking quark masses transform as a flavor-SU(3) singlet plus an octet piece:

$$\mathcal{H}_{\text{SB}} = \hat{m}(\bar{u}u + \bar{d}d) + m_s \bar{s}s = c_0 u_0 + c_8 u_8 \ , \tag{4}$$

where u_0 and u_8 are the singlet and octet scalar densities, respectively. Thus, $c_8 = (1/\sqrt{3})(\hat{m} - m_s)$ and $u_8 = (1/\sqrt{3})(\bar{u}u + \bar{d}d - 2\bar{s}s)$, etc. The matrix element $\langle N \mid c_8 u_8 \mid N \rangle$ then represents the nucleon mass shift when the SU(3) breaking is turned on. It can be related, in the leading approximation, to the baryon mass difference:

$$\tfrac{1}{3}(\hat{m} - m_s)\langle N \mid \bar{u}u + \bar{d}d - 2\bar{s}s \mid N \rangle = M_\Lambda - M_\Xi \ , \tag{5}$$

which is about -200 MeV. The one-loop chiral perturbation corrections have been estimated to be fairly significant (≈ -50 MeV) (Ref. 7). [As an approximate and compact way to include such higher-order corrections we shall use the quadratic masses on the right-hand side: $(M_\Lambda^2 - M_\Xi^2)/2M_N \simeq -254$ MeV.] A simple comparison with Eq. (1) shows that one can rewrite the left-hand side of Eq. (5) in terms of the σ term and the ratio

$$y = \frac{\langle N \mid \bar{s}s \mid N \rangle}{\frac{1}{2}\langle N \mid \bar{u}u + \bar{d}d \mid N \rangle} \ , \tag{6}$$

which expresses the strange-quark content of the nucleon:

$$\frac{1}{3}\left[1 - \frac{m_s}{m}\right](1-y)\sigma_{\pi N} = -254 \text{ MeV} \ . \tag{7}$$

Since the nucleon has only nonstrange valence quarks, the Okubo-Zweig-Iizuka (OZI) rule[10] would lead one to expect a vanishingly small $y \approx 0$. Then, using the canonical quark-mass ratio $(m_s/\hat{m}) \simeq 25$, one would predict a value for the σ term of

$$\sigma_{\pi N}^{(0)} \simeq 32 \text{ MeV} , \tag{8}$$

which is a factor of 2 too small compared to the experimental value of Eq. (3). In fact, if the extraction from the scattering data is not in error, the 60-MeV σ term implies a significant deviation of the y ratio from zero:

$$y = 1 - \frac{\sigma_{\pi N}^{(0)}}{\sigma_{\pi N}} \simeq 0.47 . \tag{9}$$

Such a large violation of the OZI rule is not necessarily in conflict with the experimental results from the deep-inelastic lepton-nucleon scatterings which do not indicate a large strange-quark content, because in the latter only the nucleon matrix elements of twist-two operators are probed. The result (9) is, nevertheless, troubling as it also implies that the nucleon would suffer a mass shift of approximately 40% in the chiral limit (see discussion below). This is counterintuitive since one would have thought that, chiral symmetry being a good one, there should not be such large effects connected with the symmetry-breaking term. Still, this mass shift and the matrix element $\langle N \mid \bar{s}s \mid N \rangle$ are theoretical properties and are not accessible to direct physical measurements. One approach one may adopt is to investigate the possibility of translating such theoretical features into more direct physical effects. In this paper I shall show that associated with this picture of the nucleon mass and a large $\bar{s}s$ nucleon matrix element is the dominance of the Higgs-boson coupling to the nucleon by strange quarks and this, in turn, results in an enhanced Higgs-boson–nucleon coupling.

II. NUCLEON MASS AND HIGGS-BOSON–NUCLEON COUPLING

Let us first review the standard ($y \simeq 0$) picture of the nucleon mass and the Higgs-boson–nucleon coupling. The most illuminating discussion on this subject has been given by Shifman, Vainshtein, and Zakharov[11] (SVZ). We shall follow their presentation and show in what ways a large-y value modifies the final picture.

The nucleon mass can be expressed in terms of the matrix element of the trace of the energy-momentum tensor $\Theta_{\mu\nu}$:

$$\begin{aligned} M_N &= \langle N \mid \Theta_{\mu\mu} \mid N \rangle \\ &= \sum_q m_q \langle N \mid \bar{q}q \mid N \rangle + \frac{\beta(\alpha_s)}{4\alpha_s} \langle N \mid G_{\mu\nu}^a G_{\mu\nu}^a \mid N \rangle , \end{aligned} \tag{10}$$

with

$$\beta(\alpha_s) = -(11 - \tfrac{2}{3} n_q) \frac{\alpha_s^2}{2\pi} + O(\alpha_s^3) \tag{11}$$

being the QCD renormalization-group β function for n_q number of quark flavors. (Whenever an explicit number is needed in this paper, we shall take $n_q = 6$.) $G_{\mu\nu}^a$, with $a = 1, 2, \ldots, 8$, is the gluon field strength. The last term in Eq. (10) is the famous triangle trace anomaly.[12]

It is observed in SVZ that one can evaluate the scalar density bilinear in the heavy-quark fields $h = c, b, t$ via the "heavy-quark expansion":[13]

$$\bar{h}h = -\frac{\alpha_s}{12\pi m_h} G_{\mu\nu}^a G_{\mu\nu}^a + O(\Lambda^3 / m_h^3) , \tag{12}$$

m_h being the heavy-quark mass and Λ the QCD scale factor (a few hundred MeV). The m_h coefficient in the trace $\Theta_{\mu\mu}$ just cancels the m_h^{-1} in the leading factor, resulting in an expression that precisely cancels the corresponding heavy-quark loop contribution in (11). Furthermore, following the standard expectation of m_u, m_d and $\langle N \mid \bar{s}s \mid N \rangle$ being extremely small (thus a negligible matrix element for the light-quark mass terms $m_u \bar{u}u + m_d \bar{d}d + m_s \bar{s}s$), SVZ obtain an expression for the nucleon mass which is principally due to the gluon contribution:

$$M_N \simeq -\frac{9\alpha_s}{8\pi} \langle N \mid G_{\mu\nu}^a G_{\mu\nu}^a \mid N \rangle . \tag{13}$$

This expression, which is manifestly chiral symmetric, succinctly represents the standard theory of the hadronic mass. It supports the valence-quark model result that each of the three valence quarks in the nucleon has a "constituent-quark mass" of about 300 MeV which is brought about mostly by the gluon interactions.

The large values of $\sigma_{\pi N}$ and y in Eqs. (3) and (9) modify this picture by predicting the importance of light-quark mass terms, particularly the strange-quark mass term,

$$\langle N \mid m_u \bar{u}u + m_d \bar{d}d + m_s \bar{s}s \mid N \rangle \simeq 410 \text{ MeV} , \tag{14}$$

and thus reduce the gluon contribution to the nucleon mass in Eq. (13) by almost a factor of 2:

$$(M_N - 410 \text{ MeV}) \simeq -\frac{9\alpha_s}{8\pi} \langle N \mid G_{\mu\nu}^a G_{\mu\nu}^a \mid N \rangle . \tag{15}$$

In terms of the valence-quark model, this means that the constituent-quark masses are not all due to gluon interactions, but also have strong current-quark mass dependences. In the following I shall show that because of the presence of $\bar{s}s$ pairs (and because of the reduced importance of the gluon component in the nucleon) the low-energy interactions between the Higgs boson and the nucleon is dominated by the strange-quark Higgs-boson coupling term.

For definiteness, I shall work explicitly with the minimal model of a single doublet of Higgs bosons. (Extension to more general cases such as the two-doublet model, one coupled to the up-type and one coupled to the down-type quarks, is straightforward. The general conclusion is not changed.) The Yukawa coupling between quarks and the Higgs scalar ϕ is parity conserving and proportional to the quark masses [with the scale set by the Higgs-boson vacuum expectation value (VEV) of $v = 2^{-1/4} G_F^{-1/2} \approx 250$ GeV]:

$$\mathcal{L}_Y = \sum_q (m_q / v) \bar{q}q\phi . \tag{16}$$

This taken between the nucleon states then yields the nucleon–Higgs-scalar coupling $g_{\phi NN}$:

$$vg_{\phi NN}=\sum_{l} m_l\langle N\mid \bar{l}l\mid N\rangle+\sum_{h} m_h\langle N\mid \bar{h}h\mid N\rangle \ . \quad (17)$$

We have grouped the light quarks $l=u,d,s$ and the heavy quarks $h=c,b,t$ separately. In the standard approach the light-quark contribution is taken to be negligibly small: m_u,m_d and $\langle N\mid \bar{s}s\mid N\rangle\simeq 0$, and the low-energy interaction between the Higgs scalar and the nucleon is dominated by the heavy-quark term: This simply reflects the importance of the gluon contribution to the nucleon mass in the standard picture. Using Eqs. (12) and (13), we have (for $y=0$)

$$vg^{(0)}_{\phi NN}\simeq\frac{2n_h}{27}M_N=n_h\times 70\ \text{MeV} \ . \quad (18)$$

Thus, in the standard approach with $n_h=3$ the Higgs-boson–nucleon coupling is proportional to the nucleon mass and is fixed to be $g^{(0)}_{\phi NN}\approx 0.84\times 10^{-3}$. Clearly, the most uncertain part of this approach is the strange-quark contribution: m_s is much larger than m_u and m_d but it is not large enough for the heavy-quark expansion to be applicable. SVZ suggest that its contribution should be unessential. On the other hand, we shall presently see that the knowledge of $\sigma_{\pi N}$ allows us to fix this strange-quark contribution: It is in fact much larger than the value a naive application of the heavy-quark expansion would suggest, and its contribution turns out to be, by far, the most important term in the Higgs-scalar coupling to the nucleon.

With the modification brought about by a large σ term, and thus also a large ratio y, the light-quark term in Eq. (17) is no longer negligibly small: we have 60 MeV from the u,d quarks, and 350 MeV from the s-quark term, while the heavy-quark contribution is reduced to 120 MeV because the gluon component of the nucleon mass is reduced [see Eq. (15)]:

$$vg_{\phi NN}\approx(60+350)\ \text{MeV}+120\ \text{MeV}=530\ \text{MeV} \ . \quad (19)$$

Thus, in the modified picture ($y\neq 0$) the strange-quark contribution dominates the Higgs-boson–nucleon coupling. In fact, it is ten times bigger than the approximately 35 MeV per flavor coming from the u,d and heavy quarks. Furthermore, the overall size of the coupling is increased by a factor of 2.5 to $g_{\phi NN}\simeq 2.1\times 10^{-3}$. In short, our central point is that the chiral-symmetric result of Eq. (18) for the Higgs-boson–nucleon coupling is drastically changed by the chiral-symmetry-breaking effects due to the strange-quark mass term.

In principle, if the Higgs scalar boson was discovered one could use reactions involving the Higgs-boson–nucleon coupling to probe the strange-quark content of the nucleon. In the meantime, the only relevant phenomenological issue is its effects on the bounds that one can derive from such experiments conducted in search of the Higgs boson. In fact, just about all the low-energy atomic and nuclear searches of the long-range interaction mediated by a light scalar boson involve this coupling. In the following I shall provide a brief review of this phenomenology. In this discussion I shall ignore the $m_\phi\gtrsim 350$ MeV bound[14] coming from the absence of the decay $K\rightarrow\pi\phi$ because its interpretation is somewhat controversial,[15] and also ignore the Linde-Weinberg bound[16] of $m_\phi\gtrsim 7.9$ GeV because the top quark is now expected[17] to be superheavy, $m_t\gtrsim 50$ GeV being just possibly in the range to spoil this bound, and because it is inapplicable to cases beyond the minimal model of a single Higgs doublet. In this connection we also note that all other searches of the Higgs boson, such as the heavy-quarkonium decay[18] $(Q\bar{Q})\rightarrow\phi\gamma$, have not up to this point yielded any useful bound on m_ϕ. Thus, it is still relevant to study the light-Higgs-boson mass limits from the low-energy nuclear and atomic physics to be discussed below.

In the early 1970 studies of x-ray transitions between large orbits of muonic atoms reported systematic discrepancies between theory (QED) and experiment.[19] Several authors pointed out that this could be accounted for by assuming an additional interaction between the muon and the nucleus mediated by a light scalar boson with $m_\phi<30$ MeV (Refs. 20 and 21). Other independent tests of this light scalar idea were carried out: A negative result in the search of the nuclear transitions $0^+\rightarrow 0^+\phi$ (with subsequent $\phi\rightarrow e^+e^-$) led to the claim of excluding $1.03\lesssim m_\phi\lesssim 18.2$ MeV (Ref. 22). Barbieri and Ericson showed that the high-precision neutron-nucleus (Pb) scattering data could be used to set a limit of $g^2_{\phi NN}\lesssim 4.3\times 10^{-10}\,(m_\phi \text{ in MeV})^4$. Combining with an estimate of $g_{\phi NN}$ from a crude fit of the then existing muonic x-ray data, a bound of $m_\phi\gtrsim 13$ MeV was claimed.[23] This limit is still being cited in the literature.[24] Subsequently, the more important developments in setting bounds on m_ϕ have been the following ones: SVZ (Ref. 11) showed, as explained above, that the Higgs-boson–nucleon coupling was fixed to be $g^{(0)2}_{\phi NN}\simeq 0.7\times 10^{-6}$. This lowered the neutron-nucleus limit down to $m^{(0)}_\phi\gtrsim 6$ MeV. The muonic atom x-ray anomaly, results have not been held up. Currently, the best experimental limit by Beltrami *et al.*[25] on the x-ray wavelength discrepancy (between theory and experiment) is $(\delta\lambda/\lambda)=(-0.2\pm 31)\times 10^{-6}$. This can be translated into a bound of $m^{(0)}_\phi\gtrsim 9$ MeV. Similarly, the initial bounds from the nuclear $0^+\rightarrow 0^+$ transitions have disappeared.[26] Currently, the most accurate data by Freedman *et al.*[27] translate into a limit of $m^{(0)}_\phi\gtrsim 9$ MeV.

Clearly, all these limits will be modified when the $g_{\phi NN}$ value increases by a factor of 2.5 over the $g^{(0)}_{\phi NN}$ value. Simple calculations show that the neutron-nucleus scattering limit of Ref. 23 increases to $m_\phi\gtrsim 10$ MeV (from 6 MeV), that the muonic atom x-ray limit of Ref. 25 increases to $m_\phi\gtrsim 11$ MeV (from 9 MeV), and that the best bound remains to be the nuclear $0^+\rightarrow 0^+$ decays of Ref. 27 with a limit of $m_\phi\gtrsim 14$ MeV (from 9 MeV).

III. DISCUSSION

The essential issue of the $\sigma_{\pi N}$ problem is that the σ term as extracted from the πN scattering data appears to be very different from the value $\sigma^{(0)}_{\pi N}$ that one gets from a picture of the nucleon consisting mainly of the up and down quarks. If these two values are not in error then we must conclude that the strange-quark content of the nucleon is large. So the questions are as follows: Are $\sigma_{\pi N} \simeq 60$ MeV and $\sigma_{\pi N}^{(0)} \simeq 32$ MeV both correct? Can we understand a large $y \simeq 0.47$? What are the physical implications of such a counterintuitive result? We shall discuss each of the points in turn.

(a) $\sigma_{\pi N} \simeq 60$ MeV. Although over the years several independent analyses[5] have all yielded $\sigma_{\pi N}$ values that are mutually compatible, the constant concern has always been that such an extrapolation procedure still contains a significant bias and the scattering data actually do not necessarily imply a large σ term.[28] Because of the difficulty in fixing the true uncertainties of a phase-shift-analysis result, it clearly will be desirable to subject the existing phases to ever-stringent consistency tests. Thus, more high-precision measurements and tests of the low-energy πN scatterings (e.g., as those performed in Ref. 29) will be very desirable. In this connection one eagerly awaits the resolution of the discrepancy between the scattering length obtained from such phase-shift analysis and that extracted from the measurement of the $2P-1S$ x-ray transition energy in the pionic hydrogen atom.[30] (The latter is also in disagreement with the πd and π^3He pionic atomic level shifts, but is apparently compatible with $\sigma_{\pi N} \simeq \sigma_{\pi N}^{(0)}$.)

(b) $\sigma_{\pi N}^{(0)} \simeq 32$ MeV. To arrive at this result two key inputs have been used: $m_s/\hat{m} \simeq 25$ and $\langle N \mid c_8 u_8 \mid N \rangle \simeq -254$ MeV of Eq. (5). Dominguez and Langacker[28] argued that the possibility of readjusting the quark-mass ratio to get a σ term of 60 MeV is most likely ruled out. (In particular, this would imply a huge violation of the nonrenormalization theorem in the K_{l3} decays.) On the other hand, Jaffe[31,32] has raised the possibility that $|\langle N \mid c_8 u_8 \mid N \rangle|$ is actually much larger than the 200 MeV of the lowest-order perturbation theory calculation. (One needs a $\langle N \mid c_8 u_8 \mid N \rangle$ of the order -500 MeV to fix the σ-term problem.) He has performed calculations first in the context of the chiral bag model,[31] and more recently (with his co-workers) in the Nambu-Jona-Lasinio model of dynamical symmetry breaking.[33] They have shown that such a picture is indeed possible in these models. In a sense here one is replacing one counterintuitive result of a large nucleon mass shift in the chiral SU(3)×SU(3) limit by another equally counterintuitive result of a large nucleon mass shift in the flavor-SU(3) limit (the success of Gell-Mann–Okubo relations notwithstanding). In this paper I have, however, accepted the validity of (a) and (b), and have concentrated on studying case (c).

(c) $y \simeq 0.47$. With respect to this result of large OZI-rule violation, there are essentially two approaches: (i) study such a theoretical possibility in model calculations of low-energy QCD, or (ii) assuming its validity, attempt to extract from it more direct physical consequences. In approach (i) the possibility of a large y has been studied by this author[34] in the context of "topological expansion" of quark graphs[35] (in this dynamical picture, the OZI-rule violation comes about through the nonplanar diagrams), by Donoghue and Nappi[36] in the context of the three-flavor Skyrme model and the bag model, and by Khatsymovsky, Khriplovich, and Zhitnitsky[37] in the context of the QCD sum rule.[38] All suggest this as a viable theoretical possibility. As for approach (ii) of extracting phenomenological consequences, the most straightforward prediction is that for the *kaon-nucleon* σ *term*:[39]

$$\sigma_{KN} = \tfrac{1}{4}(\hat{m}+m_s)\langle N \mid \bar{u}u + \bar{d}d + 2\bar{s}s \mid N \rangle \simeq 570 \text{ MeV} ,$$

to be compared to the $y=0$ prediction of $\sigma_{KN}^{(0)} \simeq 210$ MeV. Unfortunately in this instance, a critical comparison between theory and experiment may be impossible: One expects large uncertainties in the relation between the KN amplitude and the σ term, and the extraction of the amplitude from physical data is plagued with difficulties because of the complicated nonanalytic structure of the amplitude in the relevant energy and momentum range. On the other hand, a very interesting suggestion has been made by Nelson and Kaplan that a large $\bar{s}s$ matrix element in the proton implies a huge s-wave attraction between nucleons and kaons which may induce a kaon condensate at the nuclear density accessible in heavy-ion collisions (with increased strangeness production as the experimental signal).[40] Finally, in this paper I have suggested another probe of the strange-quark content of the nucleon in the form of an enhanced Higgs boson coupling to the nucleon. [In this connection, it is also interesting to note that several authors have concluded that the recent European Muon Collaboration (EMC) data on the spin-dependent proton structure function suggest a significant contribution to the proton spin by the strange-quark sea.[41]]

It is clear that none of the above-mentioned probes are very satisfactory in resolving the $\sigma_{\pi N}$ puzzle (although the Higgs-boson coupling result has immediate relevance in our search for this elusive particle). It may well be that we have to wait for a breakthrough in lattice QCD calculations[42] which are, in principle, ideally suited to compute both the matrix elements $\langle N \mid \bar{u}u + \bar{d}d \mid N \rangle$ and $\langle N \mid \bar{s}s \mid N \rangle$, thus providing a resolution of this long-standing problem in the study of chiral symmetry and hadron structures.

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