Proton Spin in Chiral Quark Models

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Abstract

The spin and flavor fractions of constituent quarks in the proton are obtained from their chiral fluctuations involving Goldstone bosons. SU(3) breaking suggested by the mass difference between the strange and up, down quarks is included, and this improves the agreement with the data markedly.

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I. Introduction

The nonrelativistic quark model (NQM) explains many of the properties of the nucleon and its excited states as originating from three valence quarks whose dynamics is motivated by quantum chromodynamics (QCD), the gauge field theory of the strong interaction. The effective degrees of freedom at low energies are dressed or constituent quarks which are expected to emerge in the spontaneous chiral symmetry breakdown of QCD that may be described by Nambu–Jona-Lasinio models \[1\]. The light quarks of QCD become dynamical quarks with mass \(m_q(p^2)\) in this process. Upon approximating the dynamical mass by \(m_q(0) \approx m_N/3\) one can introduce the concept of a constituent quark (of the NQM) at low momentum \(p\). Along with dynamical quarks Goldstone bosons \[2\] occur as effective degrees of freedom in QCD below the chiral symmetry scale \(4\pi f_\pi \approx 1169\) MeV for \(f_\pi = 93\) MeV. Other degrees of freedom, such as gluons, are integrated out.

Chiral quark models which include these effective degrees of freedom have been developed for a long time starting with the Gell-Mann–Levy \(\sigma\) model \[3\]. The nonlinear \(\sigma\) model is a starting point for soliton or Skyrme models of the nucleon \[4\]. The latter became widely appreciated when Witten \[5\] linked the Skyrme model to the large \(N_c\) limit of QCD. Chiral bag models started with ref. \[6\] but further significant development stalled when it was recognized that their failure to treat quark and hadron boosts adequately along with the violation of translation invariance is rather difficult to correct systematically. Since dynamical quarks, and constituent quarks as their low momentum limit, became more widely accepted as appropriate degrees of freedom with growing support from NJL models, chiral quark models came to dominate the literature. \[7\]

Chiral fluctuations \(q_\uparrow, q_\downarrow \rightarrow q_\downarrow, q_\uparrow + (q\bar{q})_0\) of quarks into pseudoscalar mesons, \((q\bar{q})_0\), of the SU(3) flavor octet of \(0^-\) Goldstone bosons, were first applied to the spin problem of the proton in ref. \[8\]. It was shown that chiral dynamics can help one understand not only the reduction of the proton spin carried by the valence quarks from \(\Delta \Sigma = 1\) in the NQM to the experimental value of about \(1/3\), but also the reduction of the axial vector coupling constant.
$g_A^{(3)}$ from the NQM value $5/3$ to about $5/4$. In addition, the violation of the Gottfried sum rule [8] which signals an isospin asymmetric quark sea in the proton became plausible. Here we wish to study the effects of SU(3) breaking which are needed to explain the remaining discrepancies of the spin and quark sea observables with the data.

Subsequently the analysis was extended to the $\eta'$ meson [10] although it is generally not regarded as a Goldstone boson. A singlet pseudoscalar coupling constant that differs from that of the octet was shown to cause the quark sea to become more flavor asymmetric. Amongst the pseudoscalar mesons the $\eta'$ is the heaviest. The Noether current of the $U_A(1)$ symmetry is the singlet axial vector current whose divergence contains the $U_A(1)$ anomaly. Despite the spontaneous breakdown of the $U_L(1) \times U_R(1)$ symmetry, no corresponding Goldstone boson seems to arise because of instanton configurations with integral topological charge. Thus the properties of the $\eta'$ meson differ significantly from those of other Goldstone bosons such as the pions, kaons and the $\eta$ meson. Nonetheless, for the sake of comparing with [10] in Sect. III we include also broken U(3) flavor results on the quark spin fractions with the $\eta'$ meson. Let us now turn to the $\eta$ meson case and its problems.

The $\eta$ meson arises as the octet Goldstone boson when the chiral $SU(3)_L \times SU(3)_R$ symmetry is spontaneously broken. Predictions from PCAC are not in good agreement with experiments, e. g. its octet Goldberger-Treiman relation is violated because it predicts a fairly large $\eta NN$ coupling constant which disagrees with the much smaller value extracted from analyses of both $p\bar{p}$ collisions [11] and recent precision data from MAMI [12] on $\eta$ photoproduction off the proton at threshold. Corrections from chiral perturbation theory are of order 30% and therefore much too small to help one understand the problem of the suppressed $\eta NN$ coupling better [13].

The paper is organized as follows. In Sect. II we describe the formalism of SU(3) breaking on the quark spin fractions and the quark sea contents of the proton. In Sect. III we numerically evaluate the quark spin distributions by considering the SU(3) breaking effect as brought about by the mass splitting between the up/down and the strange quarks. The
II. SU(3) BREAKING

If the spontaneous chiral symmetry breakdown in the infrared regime of QCD is governed by chiral \( SU(3)_L \times SU(3)_R \) transformations then the effective interaction between the octet of Goldstone boson fields \( \Phi_i \) and quarks is a flavor scalar given by

\[
L_{\text{int}} = -\frac{g_A}{2f_\pi} \sum_{i=1}^{8} \bar{q} \gamma_\mu \gamma_5 \lambda_i \Phi_i q. \tag{1}
\]

This interaction will flip the polarization of the quark: \( q_\downarrow \rightarrow q_\uparrow + GB \), etc. Here \( \lambda_i, \) \( (i = 1, 2, ..., 8) \) are the Gell-Mann SU(3) flavor matrices, and \( g_A \) is the dimensionless axial vector quark coupling constant that is taken to be 1, while

\[
g_A^{(3)} = \Delta u - \Delta d = \Delta_3 = F + D = (G_A/G_V)_{n\rightarrow p}, \tag{2}
\]

is the isotriplet axial vector coupling constant of the weak decay of the neutron, and \( \Delta u, \Delta d \) and \( \Delta s \) stand for the fraction of proton spin carried by the u, d and s quarks, respectively. They are defined by the following matrix elements of the axial vector currents for the nucleon state

\[
\begin{align*}
\langle N | \bar{q}_3 \gamma_\mu \gamma_5 \frac{\lambda^3}{2} q | N \rangle &= g_A^{(3)} \bar{U}_N \gamma_\mu \gamma_5 \frac{\tau^3}{2} U_N, & g_A^{(3)} &= \Delta u - \Delta d, \tag{3} \\
\langle N | \bar{q}_8 \gamma_\mu \gamma_5 \frac{\lambda^8}{2} q | N \rangle &= g_A^{(8)} \bar{U}_N \gamma_\mu \gamma_5 U_N, & g_A^{(8)} &= \frac{1}{\sqrt{3}}(\Delta u + \Delta d - 2\Delta s), \tag{4} \\
\langle N | \bar{q}_0 \gamma_\mu \gamma_5 \frac{\lambda^0}{2} q | N \rangle &= g_A^{(0)} \bar{U}_N \gamma_\mu \gamma_5 U_N, & g_A^{(0)} &= \sqrt{\frac{2}{3}}(\Delta u + \Delta d + \Delta s), \tag{5} \\
\langle N | \bar{s} \gamma_\mu \gamma_5 s | N \rangle &= \Delta s \bar{U}_N \gamma_\mu \gamma_5 U_N. \tag{6}
\end{align*}
\]

Here \( U_N \) is the Dirac spinor of the nucleon and \( g_A^{(i)} \) for \( i = 3, 8, 0 \) are the nucleon’s isovector, hypercharge and singlet axial vector couplings, respectively. It is also common to define the hypercharge spin fraction \( \Delta_8 \) and the total spin \( \Delta \Sigma \) as

\[
\Delta_8 = \Delta u + \Delta d - 2\Delta s = 3F - D, \quad \Delta \Sigma = \Delta u + \Delta d + \Delta s. \tag{7}
\]
The SU(3) symmetric chiral quark model [8,10] that invokes Goldstone boson (off-mass-shell space-like) fluctuations of constituent valence quarks inside hadrons explains several, but not all, spin and sea quark observables of the proton. Clearly, the data [14,15] call for SU(3) breaking because some of the spin fractions such as $\Delta_3/\Delta_8 = 5/3$ and the weak axial vector coupling constant of the nucleon, $g_A^{(3)} = F + D = 0.85$ [8] and 1.12 [10], respectively, still disagree with experiments in the SU(3) symmetric case. The success of hadronic mass relations suggests that a chiral interaction which breaks the SU(3) flavor symmetry also be governed by $\lambda_8$, as it is expected to originate from the mass difference between the strange and up and down quarks (and the corresponding mass differences of the Goldstone bosons).

Writing only the flavor dependence of these interactions we therefore extend the SU(3) symmetric Eq. 1 to

$$L_{int} = \frac{g_8}{\sqrt{2}} \sum_{i=1}^{8} \bar{q}(1 + \epsilon \lambda_8) \lambda_i \Phi_i q,$$  \hspace{1cm} (8)

$$\frac{1}{\sqrt{2}} \sum_{i=1}^{8} \lambda_i \Phi_i = \begin{pmatrix}
\frac{1}{\sqrt{2}} \pi^0 + \frac{1}{\sqrt{6}} \eta & \pi^+ & K^+ \\
\pi^- & -\frac{1}{\sqrt{2}} \pi^0 + \frac{1}{\sqrt{6}} \eta & K^0 \\
K^- & \bar{K}^0 & -\frac{2}{\sqrt{6}} \eta
\end{pmatrix}.$$  \hspace{1cm} (9)

Here $g_8^2 := a \sim f^2_{\pi NN}/4\pi \approx 0.08$ where $f_{\pi NN} := g_{\pi NN} m_\pi / 2m_N$ denotes the pseudovector $\pi N$ coupling constant and $g_{\pi NN}$ the pseudoscalar one. The latter can be related to Eq. 4 via the pion’s Goldberger-Treiman relation $g_{\pi NN}/m_N = g_A^{(3)}/f_\pi$. Despite the nonperturbative nature of the chiral symmetry breakdown the interaction between quarks and Goldstone bosons is small enough for a perturbative expansion in $g_8$ to apply. Note also that $\epsilon$ is the SU(3) breaking parameter which is expected to satisfy $|\epsilon| < 1$ in line with the small constituent quark mass ratio $m_q/m_s \approx 0.5$ to 0.6.

From Eq. 8 the following transition probabilities $P(u^+ \rightarrow \pi^+ + d_\downarrow),...$ for chiral fluctuations of quarks can be organized as coefficients in the symbolic reactions:
\[ u^\uparrow \rightarrow a(1 + \frac{\epsilon}{\sqrt{3}})^2(\pi^+ + d^\downarrow) + a(1 + \frac{\epsilon}{\sqrt{3}})^2\frac{1}{6}(\eta + u^\downarrow) + a(1 + \frac{\epsilon}{\sqrt{3}})^2\frac{1}{2}(\pi^0 + u^\downarrow) + a(1 - \frac{2\epsilon}{\sqrt{3}})^2(K^+ + s^\downarrow), \]

\[ d^\uparrow \rightarrow a(-1 - \frac{\epsilon}{\sqrt{3}})^2(\pi^- + u^\downarrow) + a(1 + \frac{\epsilon}{\sqrt{3}})^2\frac{1}{6}(\eta + d^\downarrow) + a(-1 - \frac{\epsilon}{\sqrt{3}})^2\frac{1}{2}(\pi^0 + d^\downarrow) + a(1 - \frac{2\epsilon}{\sqrt{3}})^2(K^0 + s^\downarrow), \]

\[ s^\uparrow \rightarrow a(1 - \frac{2\epsilon}{\sqrt{3}})^2\frac{2}{3}(\eta + s^\downarrow) + a(-1 + \frac{2\epsilon}{\sqrt{3}})^2(K^- + u^\downarrow) + a(-1 + \frac{2\epsilon}{\sqrt{3}})^2(K^0 + d^\downarrow), \]

(10)

and similar ones for the other quark polarization. The Goldstone bosons have the usual quark composition, viz.

\[ |\pi^0\rangle = \frac{1}{\sqrt{2}}(\bar{u}u - \bar{d}d), \quad |\eta\rangle = \frac{1}{\sqrt{6}}(\bar{u}u + \bar{d}d - 2\bar{s}s), \quad |K^+\rangle = u\bar{s}, \quad \text{etc.} \]

(11)

From the u and d quark lines in Eq. (10) the total meson emission probability \( P \) of the proton is given to first order in the Goldstone fluctuations by

\[ P = a[\frac{5}{3}(1 + \frac{\epsilon}{\sqrt{3}})^2 + (1 - \frac{2\epsilon}{\sqrt{3}})^2]. \]

(12)

The polarized quark probabilities may now be read off the proton composition expression

\[ (1 - P)(\frac{5}{3}u^\uparrow + \frac{1}{3}u^\downarrow + \frac{1}{3}d^\uparrow + \frac{2}{3}d^\downarrow) + \frac{5}{3}P(u^\uparrow) + \frac{1}{3}P(u^\downarrow) + \frac{1}{3}P(d^\uparrow) + \frac{2}{3}P(d^\downarrow). \]

(13)

Since the antiquarks from Goldstone bosons are unpolarized we use \( \bar{u}^\uparrow = \bar{u}^\downarrow \) in the spin fractions \( \Delta u = u^\uparrow - u^\downarrow + \bar{u}^\uparrow - \bar{u}^\downarrow, \) etc and \( \Delta s = \Delta s_{sea}, \) \( \Delta \bar{u} = \Delta \bar{d} = \Delta \bar{s} = 0. \) Moreover, the valence quark fractions are (see the NQM values in Table 1) \( \Delta u_v = 4/3, \) \( \Delta d_v = -1/3, \) \( \Delta s_v = 0. \) Altogether then Eq. (13), in conjunction with the probabilities displayed in Eq. (10), yields the following spin fractions

\[ \Delta u = u^\uparrow - u^\downarrow = \frac{4}{3}(1 - P) - \frac{5}{9}a(1 + \frac{\epsilon}{\sqrt{3}})^2, \]

(14)

\[ \Delta d = -\frac{1}{3}(1 - P) - \frac{10}{9}a(1 + \frac{\epsilon}{\sqrt{3}})^2, \]

(15)
\[ \Delta s = -a(1 - \frac{2\epsilon}{\sqrt{3}})^2. \] (16)

If the antisymmetrization of the up and down sea quarks with the valence quarks is ignored we may assume that \( u_v = 2, \ d_v = 1, \ s_v = 0 \) and \( u_{sea} = \bar{u}, \) etc so that

\[ u = 2 + \bar{u}, \quad d = 1 + \bar{d}, \quad s = \bar{s}, \] (17)
reflecting equal sea quark and antiquark numbers. From Eqs. 10, 11, 13 we now obtain the antiquark fractions

\[ \bar{u} = 2a(1 + \frac{\epsilon}{\sqrt{3}})^2, \] (18)

\[ \bar{d} = \frac{8}{3}a(1 + \frac{\epsilon}{\sqrt{3}})^2, \] (19)

\[ \bar{s} = 3a(1 - \frac{2\epsilon}{\sqrt{3}})^2 + 3a[-\frac{1}{3}(1 + \frac{\epsilon}{\sqrt{3}})]^2. \] (20)

From Eqs. 18, 19, 20 it is obvious that the sea violates the SU(3) flavor and isospin symmetries. We also see that for the broken SU(3) case \( \bar{u}/\bar{d} = 3/4 \) is still the same as in the SU(3) symmetric case \( \epsilon = 0. \) [8]

The Gottfried sum rule

\[ I_G = \int_0^1 \frac{dx}{x} [F_2^p(x) - F_2^n(x)] = \frac{1}{3} + \frac{2}{3}(\bar{u} - \bar{d}), \] (21)

where \( x \) is the Bjorken scaling variable and \( F_{2}^{p,n}(x) \) the unpolarized nucleon structure functions, measures the isospin asymmetry, \( \bar{u} - \bar{d}, \) of the antiquarks. The antiquark flavor fractions are generally defined as

\[ f_q = (q + \bar{q})/ \sum_{q=u,d,s} (q + \bar{q}), \quad \text{for} \quad q = u, d, s, \] (22)

\[ f_3 = f_u - f_d, \quad f_8 = f_u + f_d - 2f_s, \quad f_s = 2\bar{s}/[3 + 2(\bar{u} + \bar{d} + \bar{s})]. \] (23)
III. NUMERICAL RESULTS

When SU(3) breaking is included that is consistent with the higher mass of the strange quark compared to the common up, down quark mass and is governed by the hypercharge generator \( \lambda_8 \), then nearly all observables agree with the data.

In fact, with SU(3) breaking that is characterized by the parameter \( \epsilon \) defined in Eq. 8, and the parameter values \( a = 0.12 \) and \( \epsilon = 0.2 \) (see the 4th column in Table 1) the spin fraction ratio \( \Delta_3/\Delta_8 \) increases from the value 5/3 of the SU(3) symmetric case for \( \epsilon = 0 \) [8,10] and the NQM to 2.12, which is much closer to the experimental value 2.09 \( \pm \) 0.13 [14]. The situation is similar for the fraction \( f_3/f_8 \) decreasing from the value 1/3 for \( \epsilon = 0 \) (and the NQM, cf. Table 1) to 0.24 for \( \epsilon = 0.2 \) close to the experimental value 0.23 \( \pm \) 0.05. A significant defect seems to remain despite SU(3) breaking in so far as the axial vector nucleon coupling constant \( g_A^{(3)} = \mathcal{F} + \mathcal{D} = 1.217 \) for \( \epsilon = 0.2 \) is below the experimental value 1.2573 \( \pm \) 0.0028 [18]. In view of missing relativistic effects, which are known to drive this quantity even lower, this discrepancy and possibly \( \bar{u}/\bar{d} = 3/4 \) are the only ones remaining in the broken SU(3) case. Overall, SU(3) breaking leads to markedly improved results for the \( SU(3)_L \times SU(3)_R \) chiral quark model.

Another description of quark spin fractions, where \( \epsilon_{SMW} \) parametrizes the suppression of kaon transitions only, has recently been given in [19]. Upon comparing our \( \Delta s = -a(1 - \frac{2\epsilon}{\sqrt{3}})^2 \) from Eq. 16 with their \( \Delta s = -a\epsilon_{SMW} \) we obtain \( \epsilon = (1 - \sqrt{\epsilon_{SMW}}) \frac{\sqrt{3}}{2} \), and using their fit values \( \epsilon_{SMW} \approx 0.5 - 0.6 \) we find the estimates

\[
0.195 \approx (1 - \sqrt{0.6}) \frac{\sqrt{3}}{2} < \epsilon < (1 - \sqrt{0.5}) \frac{\sqrt{3}}{2} \approx 0.25,
\]

which are in reasonable agreement with the value, \( \epsilon = 0.2 \), that we establish in Table 1.

Let us also compare with the case where the singlet \( \eta' \) meson is included in chiral meson-

\[1\] The experimental value for the antiquark fraction \( f_3/f_8 \) is obtained from the measured octet baryon masses.
quark interactions with a relative coupling constant $\zeta$ that differs from that of the octet. [10]. Despite varying the additional parameter $\zeta$, the fit in the fifth column of Table 1 for the case with SU(3) breaking hardly improves the case without $\eta'$ meson in the 4th column, except possibly for $\bar{u}/\bar{d}$ decreasing from 3/4 to 0.686. In particular, the inclusion of the $\eta'$ meson does not resolve the discrepancy with the nucleon axial vector coupling constant. Since the relative $\eta'$ coupling, $\zeta = -0.3$, turns out to be much smaller than in the SU(3) symmetric case, where $\zeta = -1.2$, the $\eta'$ meson becomes almost negligible in the broken SU(3) case.

Can the remaining discrepancies in the broken SU(3) case be better understood? As we mentioned in the introduction, the Goldberger-Treiman relation of the $\eta$ meson is in conflict with experiments which can be avoided if it couples only to the strange, but not the u and d, quarks. If we assume that to be the case in the last column of Table 1, we see that both remaining discrepancies become less pronounced. In fact, $g_A^{(3)} = F + D$ increases to 1.335 (and this high value is likely to be beneficial when relativistic effects are included) and $\bar{u}/\bar{d}$ decreases from 3/4 to 7/11, while the other spin and flavor fractions change, but not by much.

IV. SUMMARY AND CONCLUSION

We have seen that in the broken SU(3) case nearly all of the nucleon’s spin observables are reproduced by the $SU(3)_L \times SU(3)_R$ chiral quark model, where the $\eta$ meson is the conventional octet Goldstone boson. The nucleon’s axial vector coupling constant $g_A^{(3)}$ may not be large enough, though, because relativistic effects are not included here which are known to drive this quantity to lower values. Including the $\eta'$ meson in the chiral dynamics does not seem to help one much to understand better the proton spin problem.

When the $\eta$ meson is taken to couple only to the strange, but not the u and d, quarks in the chiral quark model, then $g_A^{(3)} = F + D$ increases and the fit improves for $\bar{u}/\bar{d}$ as well.

The remarkable improvement in the spin and flavor fractions from SU(3) breaking shows that such chiral quark models provide a sound phenomenological framework for understand-
ing the spin problem of the proton.

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### Table 1  Quark Spin and Sea Observables of the Proton

| Observable | Data       | NQM | $a = 0.12$ | $a = 0.12$ | $a = 0.16$ |
|------------|------------|-----|------------|------------|------------|
|            | Ref. [14]  | $\epsilon = 0.2$ | $\epsilon = 0.2$ | $\epsilon = 0.2$ | $\eta$ mod. |
| $\Delta u$ | $0.84\pm0.05$ | $4/3$ | 0.824 | 0.81 | 0.87 |
| $\Delta d$ | $-0.43\pm0.05$ | $-1/3$ | -0.39 | -0.39 | -0.47 |
| $\Delta s$ | $-0.08\pm0.05$ | 0 | -0.07 | -0.07 | -0.095 |
| $\Delta \Sigma$ | $0.30\pm0.06$ | 1 | 0.36 | 0.35 | 0.31 |
| $\Delta_3/\Delta_8$ | $2.09\pm0.13$ | $5/3$ | 2.12 | 2.13 | 2.255 |
| $F + D$ | $1.2573\pm0.0028$ | $5/3$ | 1.217 | 1.205 | 1.335 |
| $F/D$ | $0.575\pm0.016$ | $2/3$ | 0.58 | 0.58 | 0.565 |
| $\bar{u}/\bar{d}$ | $0.51\pm0.09$ | 1 | 0.75 | 0.686 | 0.636 |
| $f_3/f_8$ | $0.23\pm0.05$ | $1/3$ | 0.24 | 0.235 | 0.165 |
| $I_G$ | $0.235\pm0.026$ | $1/3$ | 0.27 | 0.25 | 0.20 |
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