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HOW CAN THE STRANGENESS CONTENT OF THE PROTON BE BOTH LARGE AND SMALL*

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ABSTRACT

Experimental data suggest that some strange quark bilinear operators have relatively large matrix elements in the proton, while others are very small. We propose a likely resolution of this problem. The mixing between the light quark $\overline{q}q$ and the $\overline{s}s$ meson states strongly depends on the spin parity of the meson multiplet. We conjecture that this pattern of OZI violation in mesons and the suppression of some strange matrix elements in the proton are caused by the same nonperturbative mechanism. We construct helicity-dependent strange quark distributions in the proton, which have the expected large x nonperturbative behavior and propose the experiments in which our hypothesis can be tested.

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1. Introduction

The issue of "strangeness" content of the proton has recently received much attention[1–8] but so far the situation remains somewhat unclear. On the one hand, the success of the non-relativistic quark model (NQM), based on the hypothesis that proton is built from u and d quarks only, might be taken as an argument indicating that the number of $\bar{s}s$ pairs in the proton ought to be small. The second argument pointing in the same direction comes from the QCD sum rule calculations of the various static and dynamic nucleon parameters (mass, magnetic moments, form factors, etc.) [9] These calculations were based on the same hypothesis and they are in good agreement with experiment. Yet another argument comes from the deep-inelastic scattering data. The mean value of the proton momentum fraction carried by strange quarks has been measured and found to be small, much less than the proton momentum fraction carried by u and d quarks. On the other hand, various theoretical arguments have been put forth to indicate that the proton contains a large number of $\bar{s}s$ pairs. The theoretical interest has been triggered the EMC measurement[10] of the polarized structure function $g_1^p(x)$. The results of this experiment can be interpreted[4,5,10] as indicating relatively large matrix element in the proton of the operator $\bar{s}\gamma_\mu\gamma_5 s$. Another experimental indication of relatively large "strange content" of the proton comes from the $\pi-N$ σ-term, from which one can infer[1,2,3,11,12] (although the theoretical analysis contains some uncertainties) that

$$\frac{\langle p|\bar{s}s|p\rangle}{\langle p|\bar{u}u+\bar{d}d+\bar{s}s|p\rangle} \sim 0.2-0.1 \qquad (1)$$

We propose here a possible resolution of this difficulty. We point out that the question "how many strange quark pairs are there in the proton" is not well defined until one specifies the operator which one wishes to discuss. In afterthought, this should not come as a surprise. In meson physics it is well known that the question of strange quark content, or the degree to which the Okubo-Zweig-Iizuka (OZI)[13] rule works, strongly depends on the spin and parity of the meson multiplet. In the pseudoscalar sector the physical η and η' mesons are close to the $SU(3)$ octet $\eta_8 \sim (\overline{u}u + \overline{d}d - 2\overline{s}s)/\sqrt{6}$ and to the singlet $\eta_1 \sim (\overline{u}u + \overline{d}d + \overline{s}s)/\sqrt{3}$, respectively. Both physical states contain a strange quarks admixture of order $\mathcal{O}(1)$. In contrast, in the vector multiplet there is the so-called "ideal mixing" of the singlet and octet, leading to a physical ϕ which is purely an $\overline{s}s$ state, while $\omega \sim (\overline{u}u + \overline{d}d)/\sqrt{2}$. A similar thing happens for spin-2 mesons f' and f, with f' being an almost pure $\overline{s}s$ state. These facts cannot be explained by perturbative QCD and must be of nonperturbative origin.

2. Instantons and OZI violation in mesons

The question of qualitative explanation for the pattern of OZI rule realization and violation in the meson multiplets was previously discussed in Ref. 14. It was suggested there that the most important nonperturbative configurations of the gluonic field in QCD vacuum are of instanton type. Consider the mixed polarization operator representing OZI rule violation

$$\Pi^{s,u+d}(x,y) = \langle 0|\, T\{j^s(x), j^u(y)+j^d(y)\}\, |0\rangle \tag{2}$$

where $j^{u,s,d}$ are vector, axial-vector, pseudoscalar, scalar and tensor currents of s,u and d quarks. Disregard perturbative gluon exchange and take quarks as moving in an external instanton field. Then the *r.h.s.* of (2) factorizes

$$\Pi^{s,u+d}(x,y) = \langle 0|\, j^s(x)\, |0\rangle \, \langle 0|\, j^u(y) + j^d(y)\, |0\rangle \tag{3}$$

and may be represented by the diagram of Fig. 1.

It was shown in Ref. 14 that the matrix elements $\langle 0|\, j(x)\, |0\rangle$ in the dilute instanton gas approximation are zero for vector and tensor currents and nonvanishing for axial vector, pseudoscalar and scalar currents. For vector current the proof is simple. Consider the transformation $G_c = CI_c$, where C is the charge conjugation and I_c is the 180^o rotation around y-axis of the instanton in the $SU(2)$ subgroup of the $SU(3)$ color group. The Lagrangian of quarks in the instanton field is invariant under G_c. But the operator $j_\mu = \overline{q}\gamma_\mu q$ changes its sign under the action of G_c. Therefore $\langle 0|\, j_\mu\, |0\rangle = 0$. One can also give a proof for tensor mesons if we consider the quark energy-momentum tensor $\theta^q_{\mu\nu}$ as their source.[14] In contrast, for axial vector $j^5_\mu = \overline{q}\gamma_\mu\gamma_5 q$, pseudoscalar $j_5 \sim \partial_\mu j^5_\mu$ and scalar currents the mixed polarization operators are nonvanishing. In such a way a qualitative explanation of the OZI rule in meson multiplets can be achieved. Unfortunately the attempts[15,16] to obtain a quantitative description are model-dependent, since the dilute instanton gas approximation is not self-consistent.

It is necessary to mention here an essential feature of the proposed mechanism: the strong momentum dependence of the polarization operator (2) in momentum space. For the axial-vector case at large Q^2

$$\Pi^{s,u+d}_{\mu\nu} \sim \frac{Q_\mu Q_\nu}{Q^{12}} \tag{4}$$

in the dilute instanton gas approximation. This property is also very desirable from the phenomenological point of view: it explains the disappearance of non-perturbative effects in the charmonium region (e.g. the small η_c decay width into light hadrons).

Of course, there is no full confidence in the results obtained in the dilute instanton gas approximation. But it is plausible that some features of the instanton model of the QCD vacuum, like the selection rules (vanishing of matrix elements of some operators) and the strong Q^2 dependence correctly capture the essential physics.

3. Quark Operator Mixing in the Proton

We now turn back to the problem of strange quarks in the proton. It is plausible that the relatively large values of certain strange quark bilinear operators are due to a nonperturbative mixing between the s and u or d quark bilinears, as indicated in the diagram of Fig. 2. Let us further assume that on the qualitative level the same nonperturbative mechanism is responsible for the operator mixing in both the proton and in the mesons. A direct consequence of this hypothesis is that $\langle p| \theta^s_{\mu\nu} |p\rangle$ does not receive contributions from nonperturbative effects. Since $\langle p| \theta^s_{\mu\nu} |p\rangle$ determines the fraction of proton momentum carried by strange quarks, it follows that the latter is due to perturbative effects only and ought to be small if one adopts the common belief that perturbative effects are small even at 1 GeV. An analogous conclusion can be obtained for nucleon magnetic moments and electromagnetic form factors.

The successful proton mass calculations in the QCD sum rule approach[9] are not affected by our hypothesis. In Ref. 17 it was shown that the instanton contribution to the main term in the sum rule vanishes if the interpolating field for the proton is given by the quark current according to the prescription in Ref. 9*

4. Proton Structure Functions

Let us turn now to the matrix element $\langle p| \bar{s}\gamma_\mu\gamma_5 s |p\rangle$ which is proportional to the integral

$$\Delta s \equiv \int_0^1 dx\, [s_+(x) - s_-(x)] \tag{5}$$

where $s_+(x)$, $s_-(x)$ are the strange quark distributions with helicities $+\frac{1}{2}(-\frac{1}{2})$, respectively. In the framework of our hypothesis $\langle p| \bar{s}\gamma_\mu\gamma_5 s |p\rangle \propto \Delta s$ is nonzero due to nonperturbative effects. Unfortunately, unlike the mixing in the pseudoscalar meson nonet, we cannot support this statement by direct calculation, even in the dilute gas approximation, because the momentum transfer in the diagram in Fig. 2 is equal to zero and then the dilute gas approximation breaks down.† If the momentum transfer q_μ in the diagram of Fig. 2 is not zero, say $Q^2 = -q^2 \gtrsim 1$ GeV^2, one could attempt to calculate the corresponding strange quark contribution to the proton axial form factor in the QCD sum rule approach. However, due to the drawbacks of the dilute gas approximation we see no reason to carry out such a calculation. For us it is enough to point out, along the lines of Sec. 2, that the nonperturbative gauge field configurations of instanton type lead to the mixing of $\bar{s}s$ and $\bar{u}u+\bar{d}d$ pairs in pseudoscalar meson nonet and to the formation of η and η'

⋆ The estimate in Ref. 17 of instanton corrections to quark condensate terms in the sum rules is incorrect, due to double counting.

† Attempts to describe the spin-dependent proton structure functions via instanton contributions were made in recent papers.[18,19]

states. It follows from the Goldberger-Treiman relation that up to $Q^2 \lesssim m_\eta^2$ the flavor-octet axial form factor is described by the contribution of the η meson

$$F^8_{A,\mu} \sim \overline{v}\left(\delta_{\mu\nu} - \frac{q_\mu q_\nu}{q^2 - m_\eta^2}\right)\gamma_\nu\gamma_5 v \tag{6}$$

where v is the proton spinor. (The singlet axial form factor differs from ((6)) due to anomaly.) The $\mathcal{O}(1)$ mixing of $\overline{s}s$ and $\overline{u}u+\overline{d}d$ pairs in form factors is evident. Even though the instanton-type consideration yield the correct physics in the meson sector, we do not have to rely on them exclusively. Instead, we shall put forth a working hypothesis that $\langle p|\,\overline{s}\gamma_\mu\gamma_5 s\,|p\rangle \neq 0$ is due to a nonperturbative mechanism and we will investigate the consequences. (In dilute instanton gas approximation only the term $\sim \overline{v}q_\mu\hat{q}\gamma_5 v$ in (6) appear. So the dilute instanton gas approximation cannot describe the full axial form factor (6).)

The nonperturbative piece of the proton axial form factor has a very strong Q^2 dependence, $G^s_A(Q^2) \sim (1/Q^2)^n$, where the superscript s indicates the s-quark contribution. In the dilute instanton gas approximation $n \geq 5$. The large-Q^2 behavior of the form factor is connected with the $x \to 1$ behavior of the corresponding structure functions through the Drell-Yan relation: $(1/Q^2)^n \to (1-x)^p$, $p = 2n-1$. Therefore for $x \to 1$ we expect that $s_+(x) - s_-(x) \sim (1-x)^p$, $p \gtrsim 10$. At small x, $s_+ - s_-$ is dominated by the A_1 exchange, the intercept of the A_1 Regge trajectory being close to zero.‡ Disregarding the perturbative contribution to $s_+(x) - s_-(x)$, we propose the parametrization:

$$s_+(x) - s_-(x) = B(1-x)^p \tag{7}$$

‡ See Ref. 27 for an additional discussion of this point.

The unpolarized strange quark distribution $s(x) = s_+(x) + s_-(x)$ can also be parametrized

$$s_+(x) + s_-(x) = \frac{A}{x}(1-x)^k + C(x) \tag{8}$$

The first term in the *r.h.s.* of (8) is the generally accepted form of the strange quark distribution. The $1/x$ factor comes from pomeron exchange at small x, the factor $(1-x)^k$, $k \approx 5$ corresponds to the quark-counting rule at $x \to 1$. The term $C(x)$ represents the nonperturbative contribution to $s_+ + s_-$. If we assume that the nonperturbative contribution to $\langle p| \theta^s_{\mu\nu} |p\rangle$ vanishes, then we must require that

$$\int_0^1 x\, C(x) dx = 0 \tag{9}$$

We shall show that under the above assumptions the relatively large value of $\Delta s \simeq -0.2$ proposed in Refs. 4,5,10 as an explanation for the "spin crisis" can be reconciled with the smallness of proton momentum fraction carried by the strange quarks[20]

$$V_2^s = \int_0^1 x\,[s_+(x) + s_-(x)]\, dx = 0.026 \pm 0.006 \tag{10}$$

For the time being let us ignore $C(x)$ in (8). Combining (7) and (8) we can write

$$\begin{aligned} s_+(x) &= \frac{1}{2}\left[\frac{A}{x}(1-x)^k + B(1-x)^p\right] \\ s_-(x) &= \frac{1}{2}\left[\frac{A}{x}(1-x)^k - B(1-x)^p\right] \end{aligned} \tag{11}$$

Since $s_+(x)$ and $s_-(x)$ are positive,

$$\frac{A}{x}(1-x)^k \geq |B(1-x)^p| \tag{12}$$

It follows from (12) that

$$|B| \leq A\frac{(p-k+1)^{p-k+1}}{(p-k)^{p-k}} \tag{13}$$

and

$$|\Delta s| \leq V_2^s \, \frac{k+1}{p+1} \, \frac{(p-k+1)^{p-k+1}}{(p-k)^{p-k}} \tag{14}$$

The inequality (14) is not restrictive for large p: for example, for $p = 10$ and $k = 5$, we have $|\Delta s| \leq 8V_2^s$ and there is no contradiction between $\Delta s \simeq -0.2$ and V_2^s in (10).

There seems to be no reason in general to prevent matrix elements in the operator expansion $\langle p| \, \bar{s}\gamma_\mu D_\nu D_{\nu_1} \ldots D_{\nu_n} s \, |p\rangle$ from getting non-zero value for nonperturbative configurations of gluonic field. We therefore expect that the unpolarized strange quark distribution $s_+(x) + s_-(x)$ will also contain nonperturbative components. We expect that at $x \to 1$ $C(x)$ in (8) has approximately the same behavior as (7), $C(x) \sim (1-x)^{p'}$, $p' \approx p$. At $x \lesssim 0.1$ we expect it to be comparable in magnitude with the first term in the *r.h.s.* of (8). It is clear that despite the inclusion of $C(x)$ in (8) the inequality (14) remains qualitatively unchanged. According to our hypothesis, the distributions of $\bar{q} = \bar{u} + \bar{d}$ light antiquarks also have the nonperturbative component $C_q(x) \approx C_s(x)$. Experimentally the ratio of the momenta carried by the strange and the $\bar{q}$ sea in the proton is about 0.4[21]. In the framework of our approach it means that the perturbative part of the $\bar{q}$ sea is enhanced in comparison with the strange sea.* Therefore the nonperturbative component is relatively more important in the strange sea than in the $\bar{q}$ sea.

* The relative enhancement of the perturbative $\bar{q}$ sea might be due to the contribution from the Z-diagrams and/or the suppression of the strange sea resulting from m_s being close to Λ_{QCD}.

5. Physical Consequences and Tests of the Proposed Picture

A direct experimental test of our hypothesis could be carried out by measuring strange particle production in polarized muon (or electron) scattering on a polarized proton. Because of (7), we expect a very strong x-dependence in such an experiment, $s_+ - s_- \sim (1-x)^p$, $p \approx 10$.

Another possibility is the observation of inclusive charm production in $\overline{\nu}_\mu + N \to \mu^+ + \mathrm{charm} + X$. This experiment measures the distribution $\overline{s}_+(x) + \overline{s}_-(x) = s_+(x) + s_-(x)$ and one can attempt to measure the nonperturbative component $C(x)$ in (8). The recently published data of the CCFR collaboration[22] on charm production in neutrino experiment give some indications in favor of our hypothesis. This experiment measured the $s(x)$ distribution and found that it falls with increasing x faster than the $\overline{q}$ distribution. The authors parametrize the sea distributions by $x\overline{s}(x) = a_s(1-x)^{\beta_s}$, $x\overline{q}(x) = a_q(1-x)^{\beta_q}$ and find $\beta_s = 10.8 \pm 1.0 \pm 0.7$, to be compared with $\beta_q = 6.9$[21]. From our point of view such a parametrization is not suitable. We can fit the data[22] using the parametrization (8) with

$$C(x) = C\frac{1}{x}(1-bx)(1-x)^{p'}, \quad p' \approx p \approx 12,\ k=5, \tag{15}$$

the same of the s and the q sea. The factor $1-bx$ in (15) is introduced so that eq. (9) will be fulfilled. The constant C is constrained by the positivity of $s_+(x)$ and $s_-(x)$.

The nonperturbative mechanism contributing to the sea-quark distribution can also contribute to the gluon distribution. For the time being, however, we have no good ideas how one could observe such a nonperturbative component of the gluon distribution.

The third possibility is the measurement of the axial form factor in elastic νp scattering. This involves taking the difference $d\sigma_{el}(\nu p) - d\sigma_{el}(\overline{\nu} p)$. In this form factor we expect two structures: the usual one, which must coincide[23] with the vector form factor at large Q^2, and a nonperturbative one, rapidly decreasing with Q^2. Therefore we expect that the proton axial form factor behaves differently from the electric or magnetic (Sachs) proton form factors. The latter two are well described by a dipole fit in the region of low and intermediate Q^2 up to $Q^2 \approx 10$ GeV^2. In contrast, we expect that such a fit of the axial form factor with just one axial meson mass will not reproduce the experimental data, and that some deviations from the dipole fit must be found.

6. Scalar Densities in the Proton

In the framework of our hypothesis the matrix element $\langle p|\, \overline{s}s \,|p\rangle$ can receive contributions from nonperturbative configurations (*e.g.*, instantons in the diagram of Fig. 2) and therefore it is not small. In this context one ought to address the question whether the relatively large value of $\langle p|\, \overline{s}s \,|p\rangle$ in eq. (1) is in contradiction with the proton mass calculations in the quark model and in the QCD sum approach. One often encounters in the literature the statement[24] that the contribution of strange quarks to the proton mass is given by

$$\delta_s M_p = m_s \,\langle p|\, \overline{s}s \,|s\rangle \tag{16}$$

where $m_s \approx 150$ MeV is the current mass of the s-quark. Using (1) and the value[3,11,12]

$$\langle p|\,\overline{u}u + \overline{d}d\,|p\rangle \approx 10 \tag{17}$$

and plugging in the σ-term[25,26]

$$\frac{1}{2}(m_u + m_d)\,\langle p|\,\overline{u}u + \overline{d}d\,|p\rangle \approx 60\,\mathrm{MeV} \tag{18}$$

we have

$$\langle p|\,\overline{s}s\,|p\rangle \approx 1 - 2 \tag{19}$$

and $\delta_s M_p \approx 150-300$ MeV. Intuitively, this seems rather large (at least, the second value). In view of this, and in order to be able to compare this result with the various theoretical calculations of the proton mass (quark model, QCD sum rules, etc.), it is useful to rethink the exact meaning eq. (16). The most straightforward interpretation seems to be the following. Let M_p^0 be the mass of the proton in the limit where the strange quark mass is taken to zero, m_s=0. Assuming that the proton mass depends linearly on m_s, we then define $\delta_s M_p = M_p - M_p^0$. The theoretical computation of this quantity is quite subtle. At present it is not known how it might be carried out in a reliable fashion. With the possible exception of the lattice approach, the essential difficulty is common to all theoretical calculations. For example, in the QCD sum rule approach, if one works directly at $m_s \sim 150$ MeV, then the contribution of m_s to the proton mass comes in through the operator product expansion and appears to be small, $\sim \alpha_s^2 m_s \,\langle 0|\,\overline{s}s\,|0\rangle$. Things are different, however, with the definition of $\delta_s M_p$ above (this point was discussed earlier in Ref. 3). The computation of M_p is based on the values of the light quark condensates,

$\langle 0|\,\overline{u}u\,|0\rangle = \langle 0|\,\overline{d}d\,|0\rangle$, taken from the experiment. An analogous calculation of M_p^0 requires as input the quark condensates in the limit of massless strange quark, $\langle 0|\,\overline{u}u\,|0\rangle_{m_s=0} = \langle 0|\,\overline{d}d\,|0\rangle_{m_s=0}$. These are unknown. It is known, however, that they condensate can depend on m_s in an essential way, through a nonperturbative mixing of $\overline{s}s$ and $\overline{u}u+\overline{d}d$ quark pairs, described by the polarization operator (2). Thus one cannot rule out the possibility that in such a formulation the QCD sum rule results would be compatible with the relatively large value of $\delta_s M_p$ above.

7. Summary

Let us summarize the main points of this paper. We propose the hypothesis that the matrix elements $\langle p'|\,\overline{s}\,O_n s\,\,|p\rangle$ of the various operators O_n between proton states strongly depend on the form of the operators O_n. This strong variation in the magnitude of the matrix elements of the various operators is due to the fact that some of them receive nonperturbative contributions, while others do not. Using dilute instanton gas to gain intuition, we expect such nonperturbative contributions for $O_n = \gamma_\mu\gamma_5,\ q_\mu\gamma_5,\ \gamma_5,\ 1$ and do not expect them for γ_μ, $\theta_{\mu\nu} = \gamma_\mu\partial_\nu + \gamma_\nu\partial_\mu$. A strikingly similar hierarchy of operator mixing is known to occur experimentally in the meson sector. We postulate that both phenomena are due to the same nonperturbative mechanism. The nonperturbative terms decrease rapidly with increasing momentum transfer $Q^2 = -(p' - p)^2$. In the distribution of strange quarks in the proton they correspond to terms $\sim (1-x)^p$, $p \sim 10$. The experimental consequences of our hypothesis can be tested experimentally along the lines of Section 5.

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FIGURE CAPTIONS

1) The diagram representing the mixed polarization operator in the instanton field. The instanton is depicted by the black point.

2) The diagram representing the nonperturbative mixing of $\overline{s}s$ and $\overline{u}u+\overline{d}d$ in the proton.

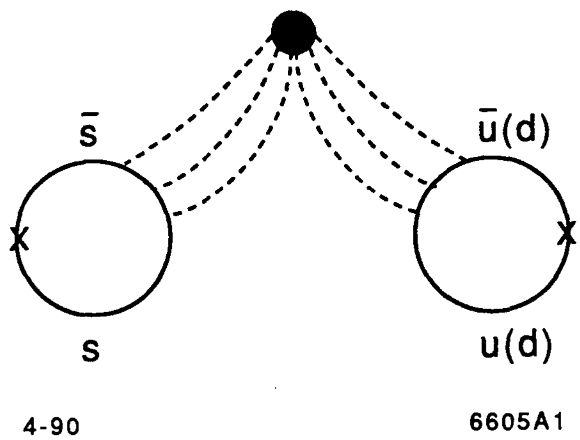

FIGURE 1

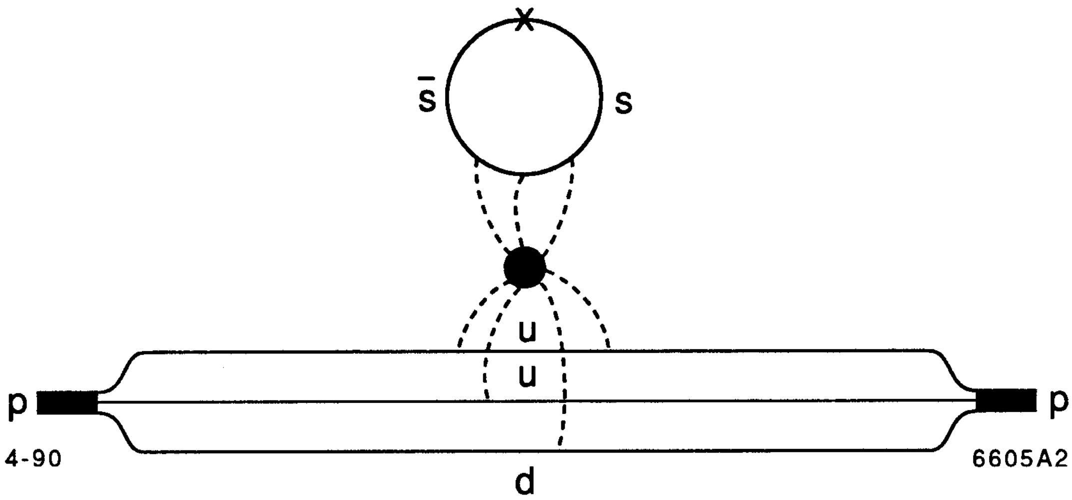

FIGURE 2