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THE LIGHT–FLAVOR STRUCTURE OF THE NUCLEON*

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Recent data on the Gottfried sum and less recent ones on the pion-nucleon sigma term seem to disagree with naive parton-based expectations on the light (*i.e.*, up, down and strange) quark content of the nucleon. We show that these discrepancies are resolved if nonperturbative contributions are included in the analysis of the data. These appear both in the computation of matrix elements of operators, and in their QCD scale dependence, and depend strongly on the quantum numbers of the given state.

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1. Introduction

In recent times a set of high-precision measurements of the nucleon structure functions, both in the polarized and unpolarized case, has led to results in striking disagreement with naive parton model expectations. The first of these to come [1] has spawned a vast theoretical literature [2], the so-called nucleon spin crisis; it is now clear [3] that the experimental result can be fully explained and understood only in a nonperturbative framework, and, more specifically, that it may be a first hint of a nonperturbative effect which survives the high-energy limit and is thus visible in the deep-inelastic regime.

More recent results on the Gottfried sum rule [4, 5] have put in doubt even the standard [6] assumptions on the up and down (unpolarized) flavor content of the nucleon. Combined with the long-standing sigma term puzzle

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[7], which hints to an unusual strange quark content, these results seem to suggest that the light-flavor structure of the nucleon is rather different from what the naive intuition based on a combination of the non-relativistic quark model with the naive parton model would suggest.

It seems thus that naive expectations have to be reconsidered. These are typically expressed as Ansätze on the form of the quark distributions $q_i^k(x)$, which in the parton model provide the probability of a quark of flavor i to carry a fraction x of the momentum of the baryon k, and are related to the nucleon structure functions. The quark distributions are split into a valence and a sea part:

$$q^b(x) = q_v^b(x) + q_s^b(x)\,; \quad \bar{q}^b(x) = q_s^b(x)\,. \tag{1}$$

In QCD this decomposition has only meaning as a definition: the quark distributions are related to the expectation value of certain operators in such a way that the structure function $F_2(x)$ is related to them by [8]:

$$F_2^k(x) = x\sum_i e_i^2\left(q_i^k(x) + \bar{q}_i^k(x)\right), \tag{2}$$

where the sum runs over all the quark flavors and e_i are electric or weak charges according to whether the target is probed with electrons or neutrinos. The quark and antiquark components for various flavors can be extracted by taking suitable linear combinations of structure functions for neutrino, antineutrino and electron scattering off different targets. QCD then predicts the scale dependence of these quantities.

However, it is usually assumed that the quark distributions can be actually viewed as expressing the quark content of the given hadron, and that this picture provides information on the nucleon structure in that it suggests the form of the distribution themselves. This assumption is based on the expectation that the matrix elements of operators which provide the quark–parton distributions (and which can only be written down in the light-cone formalism [9]) can be identified, at least up to some unknown renormalization, with the quark operators of the naive quark model. Such an assumption is totally unjustified within QCD; it is however quite successful phenomenologically, and it is customarily used in analyzing the experimental data.

A simple example of application of this wisdom is the computation of the ratio $R(x) = F_2^n(x)/F_2^p(x)$ as $x \to 1$. The naive argument goes as follows: as $x \to 1$ only the valence quarks survive, because if a quark is carrying the whole of the nucleon's momentum it must be a valence one. Furthermore, if one single quark is carrying the quantum numbers of a nucleon, by isospin, it ought to be a u quark for the proton and a d quark for the neutron. Thus

$\lim_{x\to 1} R(x) = e_d^2/e_u^2 = 1/4$. This prediction is spectacularly borne out by the data [4, 10], while it is clearly based partly on symmetry (isospin) and partly on naive guesswork.

More in general, the quark distributions (1) are usually taken to satisfy the following assumptions:

1. The sea and valence components separately satisfy isospin symmetry:

$$u_v^p = d_v^n, \quad u_v^n = d_v^p; \qquad u_s^p = d_s^n, \quad u_s^n = d_s^p. \tag{3}$$

2. The sea component in each baryon is SU(2) symmetric:

$$u_s^p = d_s^p; \quad u_s^n = d_s^n. \tag{4}$$

3. The content of Zweig suppressed flavors, such as strange quarks in the nucleon, is small. This implies relations like

$$\langle p|\bar{s}s|p\rangle \approx \langle n|\bar{s}s|n\rangle \approx 0. \tag{5}$$

It turns out that the recent data on the Gottfried sum rule put either (or both) assumptions 1 and 2 into question, and the long-standing sigma term puzzle challenges assumption 3. In the sequel, we will discuss both these points in turn. We shall show that the naive picture can actually be recovered, provided nonperturbative QCD effects are taken into account, both in the QCD evolution and in the determination of the symmetry structure of matrix elements of operators.

2. The puzzle of the Gottfried sum

The question of the validity of assumptions 1 and 2 can be condensed into an analysis of the Gottfried sum S_G, defined by

$$S_G \equiv \int dx \, \frac{F_2^p(x) - F_2^n(x)}{x}. \tag{6}$$

According to Eq. (2), the first moment of F_2 corresponds to the sum of the electric charges squared of all partons present in the target,

$$\int dx \, \frac{F_2(x)}{x} = \sum_i e_i^2 \int dx \, (q_i(x) + \bar{q}_i(x)) \equiv \sum_i e_i^2 n_i, \tag{7}$$

where we have defined

$$n_i = \int dx \, (q_i(x) + \bar{q}_i(x)). \tag{8}$$

This is expected to be a divergent quantity for a single nucleon, due to the fact that at small x the nucleon's momentum is subdivided among an infinite number of quarks; however this divergence (*i.e.*, the small-x behavior of all $q^i(x)$) is expected to be universal (on the basis of Regge theory) [11], essentially due to the flavor-independence of the strong interaction, hence S_G Eq. (6) is expected to be finite.

The symmetry assumptions (1), (2) amount to the further conjecture that not only in the $x \to 0$ limit, but actually for all x the sea is universal, hence only valence quarks can contribute to nonsinglet combinations such as Eq. (6). Taking only the valence component into account one arrives thus at the so-called Gottfried sum rule [11, 12]:

$$S_G = \tfrac{4}{9} n^p_{u,v} + \tfrac{1}{9} n^p_{d,v} - \tfrac{4}{9} n^n_{u,v} - \tfrac{1}{9} n^n_{d,v} \tag{9}$$

$$= \left(\tfrac{4}{9} - \tfrac{1}{9}\right)\left(n^p_{u,v} - n^p_{d,v}\right) = \tfrac{1}{3}\,. \tag{10}$$

However S_G has now been measured [4] as:

$$S_G = 0.24 \pm 0.016 \text{ at } 4 \text{ GeV}^2\,. \tag{11}$$

Even though the extraction of this value involves certain assumptions, *e.g.* an extrapolation of the measured structure function to the full range of x, it seems exceedingly unlikely that a value compatible with the naive expectation Eq. (10) could be obtained. As a matter of fact, an improved data analysis [5] has given the new value

$$S_G = 0.258 \pm 0.017 \text{ at } 4 \text{ GeV}^2 \tag{12}$$

but has also shown evidence for shadowing in deuterium. If shadowing effects (which cannot be measured precisely as yet) were included, the value of S_G as currently determined (Eq. (11) or (12) should be brought down by an amount which is theoretically estimated [13] to be $\Delta S_G \sim 0.02 - 0.04$. The value Eq. (11) seems thus a conservative estimate of the value of S_G, which might turn out to be even lower.

It is thus necessary to reconsider assumptions 1 and 2. If we make no assumption, the Gottfried sum is generally given by

$$S_G = \tfrac{1}{2}\left[\tfrac{5}{9}(n_u + n_d) + \tfrac{1}{3}(n_u - n_d)\right]_{I=1}, \tag{13}$$

where the subscript indicates that the isotriplet part must be taken, *i.e.* the difference of the expression in square brackets between a proton and a neutron. Assumption 1 (isospin symmetry) implies $[n_u + n_d]_{I=1} = 0$, whereas assumption 2 (sea flavor symmetry) implies that $[n_u - n_d]_{I=1}$ be identified with its valence value, *i.e.*, $[n_u - n_d]_{I=1} = 1$. The data could be explained either by assuming that the neutron sea is bigger than the proton one (thus violating isospin), or that the proton sea and separately the neutron sea are not flavor neutral, but rather they have a flavor asymmetry that anticorrelates to that of the valence component (thus violating sea flavor symmetry).

In order to fix the relative amount of violation of these two symmetries one needs an independent measurement of a different linear combinations of the quantities entering Eq. (13). It can be shown [14, 15] that the required information is provided by the isotriplet part σ_G of the so-called nucleon σ-term, defined as

$$\sigma_G \equiv \sigma_{I=1} = \langle N|m_u\bar{u}u + m_d\bar{d}d|N\rangle_{I=1}\,, \tag{14}$$

where m_i are current quark masses. This is because σ_G satisfies [14, 15]

$$\sigma_G = \tfrac{1}{2}\left[(m_u + m_d)(n_u + n_d) + (m_u - m_d)(n_u - n_d)\right]_{I=1} = M_p - M_n \tag{15}$$

which is the analogue of the Gottfried sum S_G, with square charges replaced by masses. Using this extra piece of information one gets to the conclusion that

$$\tfrac{1}{2}[n_u + n_d]_{I=1} \approx \pm 0.03 \tag{16}$$

$$\tfrac{1}{2}[n_u - n_d]_{I=1} - 1 \approx -0.35, \pm 0.1 \tag{17}$$

i.e., sea flavor neutrality is broken considerably more strongly than isospin.

Just as in the well-know case of the "spin puzzle" [1, 2] it turns out that the experimental result Eqs (16), (17) can be easily accommodated in effective models of the nucleon, such as the Skyrme or the bag model [14]. Thus, again like in the spin case, one may ask whether the data are not simply telling us that the "naive parton wisdom" is simply not applicable in this channel.

To make this more precise, one should first however investigate what QCD does tell us. Any nonsinglet combination of first moments of structure functions $F_2(x)$, such as S_G Eq. (6) acquires a two-loop scale dependence via the Altarelli-Parisi equations [8]. Indeed, it is clear that at one loop the quark number does not evolve, because at one loop the only elementary process (in the nonsinglet channel) is gluon radiation by a quark line (Fig. 1), which clearly leaves the total number of quarks (albeit not their momentum distribution) unchanged. At two loops, however, the emitted gluon can in turn produce a quark-antiquark pair through the diagram of Fig. 2, thereby increasing the total quark plus antiquark number measured by the first moment of $F_2(x)$. At first sight it would seem that this again should not contribute to the evolution of nonsinglet quantities, because the probability for pair emission is independent of the emitted flavor. However, this is not entirely correct because if the emitted pair has the same flavor as the original quark the final state must be antisymmetrized with respect to the two identical quarks, while this is of course not the case if the flavor of the radiated pair is different [16]. This is enough to generate a tiny difference between the splitting function $\mathcal{P}^{D}_{qq}$ for flavor-diagonal emission (*i.e.*, the case when the emitted pair has the same flavor as the original quark) and the splitting function $\mathcal{P}^{ND}_{qq}$ for flavor-nondiagonal emission (*i.e.*, the case when the emitted pair has a different flavor).

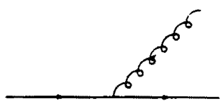

Fig. 1. The gluon radiation diagram which leads to one-loop Altarelli–Parisi evolution.

It follows that S_{G} acquires a scale dependence [16]; this is given by the two-loop Altarelli-Parisi equation [8]

$$\frac{d}{dt}(q \pm \bar{q})^{\mathrm{NS}} = (q \pm \bar{q})^{\mathrm{NS}} \otimes (Q_{qq} \pm Q_{q\bar{q}}) \ , \qquad (18)$$

where $\otimes$ indicated convolution of the splitting functions and the quark distributions, NS indicates any nonsinglet combination of quark distributions, and the splitting function responsible for the nonsinglet evolution is given by

$$Q = \mathcal{P}^{\mathrm{D}} - \mathcal{P}^{\mathrm{ND}}$$

in terms of the diagonal and nondiagonal splitting functions. It follows that the evolution of S_{G} Eq. (6) is given by

$$\frac{d}{dt} S_{\mathrm{G}}(t) = \left(Q^{1}_{qq} + Q^{1}_{q\bar{q}}\right) S_{\mathrm{G}}(t), \qquad (20)$$

in terms of the first moments Q^{1}_{qq}, $Q^{1}_{q\bar{q}}$ of the various splitting functions. The physical interpretation of Eqs (18), (20) (which is derived *e.g.* in Ref. [8]) is straightforward: the nonsinglet evolution is due to the difference in probability for flavor-diagonal and flavor-nondiagonal pair production; if the latter is greater than the former then the nonsinglet quark number measured by S_{G} decreases because the sea generated from the process of Fig. 2 tends to compensate the starting flavor asymmetry.

The same results can be obtained [16] via operator-product expansion methods, which are however much more cumbersome because only even moments of the (charge-conjugation even) combination of structure functions

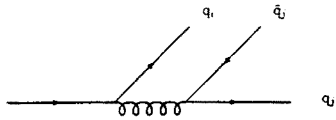

Fig. 2. The two-loop diagram which generates nonsinglet evolution of quark distribution. When $i = j$ the final state must be antisymmetrized with respect to the two identical quarks.

which appears in Eq. (6) have a direct interpretations in terms of matrix elements of twist-two local operators, whereas the odd moments (such as the first moment which defines S_G) have to be constructed by analytic continuation. Anyway, S_G is thus found to evolve perturbatively according to

$$S_G(Q^2) = \left[1 + \frac{\gamma^{(2)}}{b}\left(\alpha_s(Q^2) - \alpha_s(\mu^2)\right)\right] S_G(\mu^2), \tag{21}$$

where explicit perturbative computation up to two loops leads to $\gamma^{(2)}/b = 0.01$ [16]. It follows that perturbative QCD evolution, even though present, is entirely negligible from a quantitative point of view because if $\alpha_s < 1$ (as it must in order for perturbation theory to make sense) then the total variation of S_G is less than 1 %.

It seems thus that, just like in the spin case, even if we assume that the "parton wisdom" should apply at some low scale, while the experimental result is taken at large Q^2, perturbative QCD evolution cannot explain the discrepancy between the experimental value Eq. (11) and the expectation Eq. (10). More specifically, even if we are willing to give up the naive expectation Eq. (10) we must accept the rather unpalatable conclusion that even at very low scales the quark-parton content of a nucleon is very different from its constituent quark content.

There is, however, a way out of this impasse, which consists of allowing for the possibility that the perturbative evolution Eq. (21) is replaced by some nonperturbative evolution at low enough scales. Due to a nonperturbative mechanism this evolution could then be much larger than the perturbative one. Symmetry considerations suggest what the origin of such a mechanism may be [17]. If the interaction which governs the generation of the quark-antiquark sea is $U(N_f)$-symmetric, then by Wigner–Eckhart's theorem the sea will be flavor symmetric. In such case there is no QCD evolution at all, and indeed Eq. (18) shows that all non-singlet scale dependence can be traced to a difference in diagonal and non-diagonal splitting functions which signals the violation of exact $U(N_f)$ symmetry. The evolution is thus, crudely speaking, a measure of the amount of $U(N_f)$ symmetry breaking. Now, it is well known that in QCD the axial U_A (N_f) flavor symmetry is dynamically broken down to a SU_A (N_f) by the axial anomaly. This is manifested in the mass spectrum of pseudoscalar mesons, since the singlet meson (the η') has a much larger mass than the octet mesons (pions and kaons) which can be viewed as Goldstone bosons. If such effect is incorporated in the Altarelli–Parisi equations it could thus potentially lead to large evolution. Notice that this can only be done in an effective way, since there is no known way of obtaining this pattern of symmetry breaking through purely perturbative computations (rather, nonperturbative effects such as instantons are thought to be at the origin of it).

We can thus try to effectively include nonperturbative symmetry-breaking effects by considering contributions to the evolution equations where a meson couples directly to the quark-partons. These can be viewed pictorially as being given by diagrams with the same topology as those of Fig. 2, but where the original quark and the emitted antiquark exchange many gluons in such a way that they can be viewed as a quark-antiquark bound states, thus leading to the diagram depicted in Fig. 3. It is clear that, qualitatively, this should give a contribution to the Altarelli–Parisi evolution which may potentially generate significant evolution with the right sign as required to explain the observed effect Eq. (11). Indeed, due to the larger mass of the η', the singlet meson emission should be energetically suppressed in comparison to octet meson emission. Because flavor-nondiagonal emission is entirely due to the latter process, whereas flavor-diagonal emission proceeds in part through the former, one expects that the diagonal process will be disfavored. If Eq. (18) may be generalized to meson emission, this should result in a negative value of its r.h.s., thus leading to a partial screening of the valence flavor asymmetry. Notice that the process given by the elementary diagram of Fig. 3 is now effectively a one-loop one, thus the corresponding evolution is potentially strong.

Fig. 3. The meson radiation diagram which generates nonperturbative evolution.

It seems thus reasonable to investigate this option in a quantitative way [18]. The required generalization of the Altarelli–Parisi evolution proceeds in two steps: first, we must generalize the set of evolution equations so as to include the coupling of quarks to bound states, and then we must compute the new splitting functions which enter the generalized equations. Because we are only interested in nonsinglet evolution, we may concentrate on the pseudoscalar meson sector, since for all other mesons or baryons the flavor symmetry will be approximately exact, thus, by the above argument, no contributions to the evolution should arise.

The main difference between the usual Altarelli-Parisi equations and the generalized ones is then the possibility of flavor-nondiagonal transitions already at the one loop level, such as that given by the diagram of Fig. 3 when one starts with, say, a u quark and ends with a d quark and a π^+. It is then easy to see that the nonsinglet evolution equations generalize (considering for simplicity flavor SU(2) only) as

$$\frac{d}{dt}q = \mathcal{Q}_{qq} \otimes q + \mathcal{Q}_{q\bar{q}} \otimes \bar{q} + \mathcal{P}_{q\pi} \otimes \pi , \tag{22}$$

$$\frac{d}{dt}\bar{q} = \mathcal{Q}_{qq} \otimes \bar{q} + \mathcal{Q}_{q\bar{q}} \otimes q - \mathcal{P}_{q\pi} \otimes \pi , \tag{23}$$

$$\frac{d}{dt}\pi = \mathcal{P}_{\pi q} \otimes (q - \bar{q}) + \mathcal{P}_{\pi\pi} \otimes \pi , \tag{24}$$

where q, $\bar{q}$ and π are quark, antiquark and pion nonsinglet distributions, respectively, defined as

$$q(x) = q_u(x) - q_d(x) \tag{25}$$

$$\bar{q}(x) = \bar{q}_u(x) - \bar{q}_d(x) \tag{26}$$

$$\pi(x;t) = \pi^+(x;t) - \pi^-(x;t) , \tag{27}$$

and, at one loop, the splitting functions $\mathcal{P}^D$ and $\mathcal{P}^{\rm ND}$ are obtained from linear combinations of the diagrams of Fig. 3 corresponding to various pseudoscalar mesons.

Combining Eqs (22)–(24) it is easy to show that that the evolution equation for quark distribution becomes

$$\frac{d}{dt}(q + \bar{q}) = (q + \bar{q}) \otimes (Q_{qq} + Q_{q\bar{q}}) , \tag{28}$$

$$\frac{d}{dt}(q - \bar{q}) = (q - \bar{q}) \otimes (Q_{qq} - Q_{q\bar{q}}) + \mathcal{P}_{q\pi} \otimes 2\pi , \tag{29}$$

where the various distributions and splitting functions are those which appear in Eqs (22)–(27). It follows that the evolution equation of the Gottfried sum Eq. (20) is formally unchanged, provided of course the $\mathcal{Q}_{qq}$ and $\mathcal{Q}_{q\bar{q}}$ are the new splitting functions which include meson emission. Notice that the meson distributions and splitting functions π, $\mathcal{P}_{\pi q}$ and $\mathcal{P}_{q\pi}$ do not directly contribute to the evolution of $q+\bar{q}$, and thus of $S_{\rm G}$, whereas they do contribute to the evolution of $q - \bar{q}$.[1] In order to avoid double counting the contributions to the QCD evolution determined in this guise should not be added on top of the usual perturbative evolution. Rather, at large enough Q^2 (presumably not larger than $Q^2 \sim 10$ GeV^2) the nonperturbative evolution should flatten out, so that at some large scale the perturbative behavior is reproduced.

The next required step is the computation of the splitting functions. Now, the whole picture of Q^2 evolution as being described by a probabilistic process relies on the dominance of certain diagrams (ladder diagrams) which contribute to the deep-inelastic scattering process, and are then effectively resummed by solving the evolution equations. This corresponds physically to the dominance of processes where the emitted particle (a gluon in the standard evolution equations, a meson in the present case) is quasi-collinear to the original one (Weizsäcker–Williams approximation) and has been proven [20] to hold for gauge theories (both abelian and nonabelian) in specific gauges in the Bjorken limit. These proofs can actually be also extended [20] to the case of theories with "pseudoscalar glue", *i.e.*, to the case of coupling of fermions to point-like pseudoscalars. It is much less clear that such a picture should hold in the case of extended pseudoscalars: as a matter of fact, coupling to pseudoscalar glue for all values of Q^2 would lead to scaling laws for the moments of structure functions which are not in agreement with the experimental data [8], and the effects discussed here can only be important in a limited Q^2 range, not so large that the quark-antiquark structure of the mesons is seen.

Hence, we proceed to the computation of the splitting function by assuming this approximation to hold. We verify then that this assumption is consistent, in that either significant Q^2 evolution occurs and the meson emission cross section is dominated by a kinematical range where the meson is effectively pointlike, or, if the non-pointlike region is providing an important contribution to the cross-section, the evolution is negligible. This must happen at a value of Q^2 low enough that the know asymptotic saling

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[1] This means that the so-called Sullivan process, deep-inelastic scattering off quarks inside pions in the nucleon wave function, cannot contribute directly to measurements of S_G. This is in keeping with the naive observation that a pion, being made of an equal number of u and d quarks (if quarks and antiquarks are counted with the same sign) cannot contribute to S_G. Pion emission may, nevertheless, contribute *indirectly* to S_G by modifying the flavor structure of the sea. Thus whereas pion counting does not tell us anything about S_G (contrary to occasional naive statements) pion emission mechanism (sometimes suggested within effective models [19]) may lead to a contribution to S_G.

properties are unaffected. The splitting functions are then given by

$$\left[\mathcal{P}_{q_i q_j}(x;t)\right]_{\Pi} = \frac{d}{dt}\sigma^{\gamma^*\Pi}_{q_i q_j}(x;t), \tag{30}$$

in terms of the total cross-section $\sigma^{\gamma^*\Pi}_{q_i q_j}(x;t)$ for the process of Fig. 3, integrated over all transverse momenta and expressed in terms of the usual Bjorken variables. Using this definition in the evolution equation (20) the evolution of $S_{\rm G}$ is immediately found to be given by the multiplicative law

$$S_{\rm G}(t) = \Delta(t,t_0)S_{\rm G}(t_0), \tag{31}$$

where the nonperturbative determination of $\Delta(t,t_0)$ is

$$[\Delta_1(t,t_0)]_{\rm nonpert} = \exp\Bigg(\left[\tfrac{1}{6}\sigma_1^{\gamma^*\eta'}(t) + \tfrac{1}{3}\sigma_1^{\gamma^*\eta'}(t) - \tfrac{1}{2}\sigma_1^{\gamma^*\pi}(t)\right] - \left[\tfrac{1}{6}\sigma_1^{\gamma^*\eta'}(t_0) + \tfrac{1}{3}\sigma_1^{\gamma^*\eta'}(t_0) - \tfrac{1}{2}\sigma_1^{\gamma^*\pi}(t_0)\right]\Bigg). \tag{32}$$

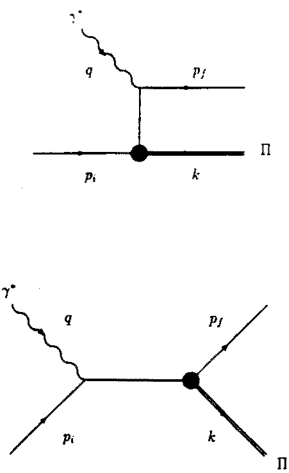


Fig. 4. Deep-inelastic scattering off a quark which radiates a bound state Π: a) t-channel diagram; b) s-channel diagram.

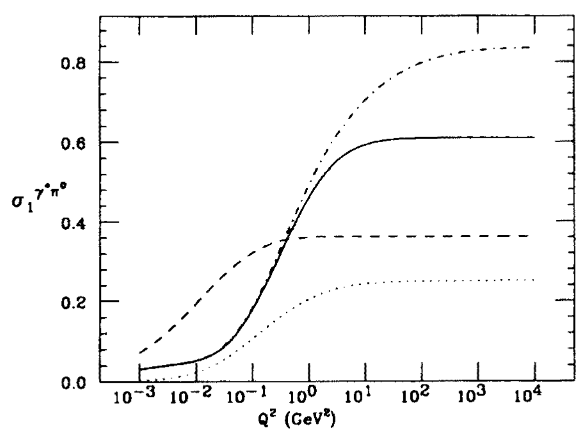
The problem of determining the evolution of $S_{\rm G}$ is thus reduced to the computation of the various cross-sections which enter Eq. (32). The computation, which is somewhat involved technically, has been performed and discussed in Ref. [18]. Because the pion-quark coupling is of course unknown, the coupling has been written in the most general form compatible with Lorentz invariance; this requires introducing four independent form factors (a pseudoscalar, two axial, and a tensor). The cross-section is then computed from the two diagrams (corresponding to s- and t-channel contributions) of Fig. 4. Using the asymptotic behavior of the form factors for large values of their arguments which is known from the behavior of the pion vertex function, it is found that for intermediate values of Q^2 (up to a few GeV^2) the cross-sections are controlled essentially by the pseudoscalar form factor, while the strength of one of the axial couplings controls the large Q^2 tail. The pseudoscalar form factor can be further constrained by use of the Ward identity and a condition which relates it to the physical value of the pion decay constant f_π. It turns out that all the evolution up to intermediate Q^2 is thus parametrized by a single parameter which can be identified with the constituent quark mass, while an extra parameter (the axial coupling) controls the large Q^2 tail.

The main qualitative features of the result of this computation can be summarized as follows:

1. The cross-sections are indeed dominated by the region where the form factor is pointlike; this is the region where the emitted meson and quark are almost collinear, consistently with the parton interpretation. This is shown in Fig. 5, where the cross-section for pion emission is separated into its various pieces (Fig. 5a), which are then displayed (Figs 5b, 5c) for various values of Q^2 as a function of x. The pointlike dominance is apparent by noting that the form factor in the s and t channel are, respectively, $\varphi_{\hat{s}} = \varphi\left(\frac{1}{4}(2\hat{s} - M^2)\right)$, and $\varphi_{\hat{t}} = \varphi\left(\frac{1}{4}(M^2 - 2\hat{t})\right)$ with

$$\varphi^\pi(p^2) = \frac{m_d}{f_\pi}\frac{\Lambda^2 + m_d^2}{\Lambda^2 + p^2}\,; \qquad (33)$$

and $s = \frac{(1-x)}{x}Q^2$ as $x \to 1$, while $t = -\frac{x}{1-x}m_\pi^2$ in the collinear $k_{\rm perp} \to 0$ limit, implying that the collinear, pointlike limit is $x \to 1$ in the



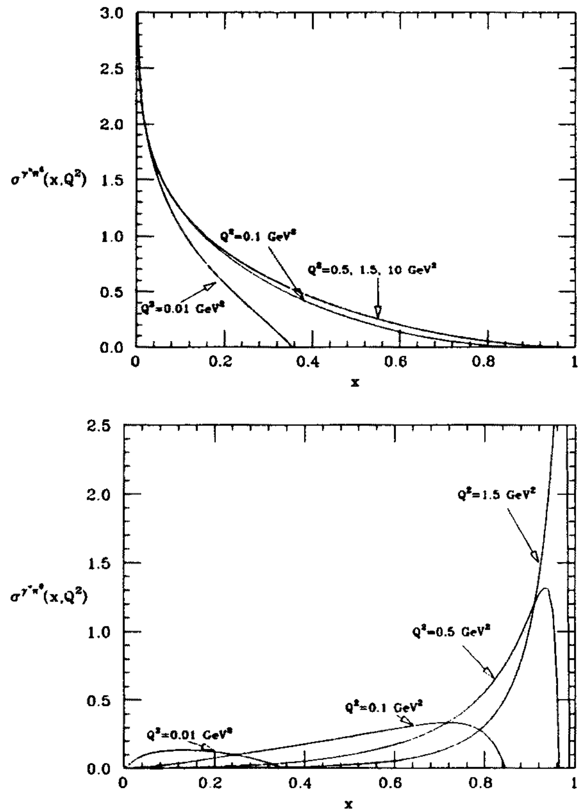


Fig. 5. The π^0 emission cross-section (with constituent quark mass $M = 325$ MeV) computed from the two diagrams of Fig. 4: a) cross section integrated over x as a function of Q^2; full line: full cross section with no axial coupling; dot-dash line: full cross section with axial coupling; dashed line: t-channel contribution (no axial coupling), diagram Fig. 4a; dotted line: s-channel contribution (no axial coupling), diagram Fig. 4b. b) t-channel contribution (dashed line of Fig. 5a) as a function of x and Q^2. c) s-channel contribution (dotted line of Fig. 5a) as a function of x and Q^2.

s channel and $x \to 0$ in the t channel. It is interesting to observe that Figs 5b and 5c show that the growth of the cross-section with Q^2 (which determines, according to Eq. (30), the anomalous dimensions) is due to the growth of the collinear peak, which dominates the cross-section for all Q^2 larger than ~ 0.1 GeV^2. In the very small Q^2 region ($Q^2 < 0.1$ GeV^2) where the non-pointlike contributions are important the Q^2 dependence of the cross-section flattens (Fig. 5a).

2. The axial coupling controls only the high-Q^2 tail of the cross-section (see Fig. 5a), while the bulk of the Q^2 dependence of the cross-section

occurs in a region ($Q^2 < 5$ GeV^2) where the cross-section is controlled by the pseudoscalar coupling.

3. The anomalous dimensions are positive definite throughout the Q^2 range. The anomalous dimension for pion emission is indeed much larger than that for eta emission, thus leading to screening of the Gottfried sum, according to Eqs (31), (32). This is displayed in Figs 6a and 6b where the anomalous dimension calculated from Eq. (30) is shown in the extreme cases of π^0 and η' emission.
4. All anomalous dimensions flatten at low Q^2 (Fig. 6), thus so does necessarily the evolution of S_G. It follows that this can be smoothly connected at the low scale where flattening occurs to a quark model predictions.

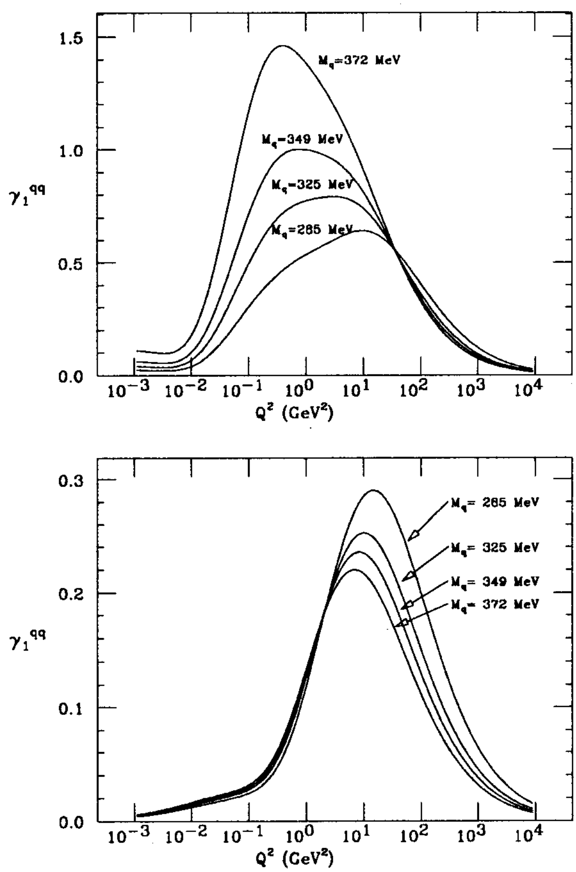


Fig. 6. Anomalous dimension calculated from Eq. (30) for meson emission, for different values of the constituent quark mass: a) π^0 emission; b) η' emission. The axial coupling is set to a large (*i.e.* maximal) value.

5. All anomalous dimensions flatten at large Q^2 (Fig.6), due to the fact that the cross-section for meson emission behaves as an inverse power of Q^2 when Q^2 is large enough. The evolution of S_G flattens more rapidly than any of the various evolutions due to meson emission separately does, since it is proportional to a difference of mesonic cross-sections, Eq.(32). It follows that the mesonic effects disappear at large Q^2, where they behave as higher-twist effects, and perturbative behavior is regained.

We proceed thus to a computation of the evolution of S_G according to Eqs (31), (32). The starting value of S_G is chosen as the naive quark model expectation Eq. (10). The scale at which this value holds is fixed either as $|Q| = 200$ MeV, or as the scale at which the evolution flattens. Interestingly, this leads to the same result (small variations of the flattening point being irrelevant due to the multiplicative nature of the evolution).

The final result for the evolution of S_G is shown in Fig.7, for different values of the constituent quark mass (Fig. 7a) and for different values of the axial coupling which controls the high-Q^2 tail (Fig. 7b). This shows that the experimental value is reproduced as a consequence of nonperturbative evolution from the naive quark model value. In the region where the data are taken or slightly above the nonperturbative evolution flattens away; our computation provides an effective smooth interpolation (which is however sensitive to a free parameter of the computation, namely, the axial coupling) between the nonperturbative evolution and the perturbative behavior which holds in the high Q^2 limit. A robust prediction of this computation is that significant evolution (characterized by anomalous dimensions of several orders of magnitude larger than the perturbative one) is taking place in the intermediate Q^2 range, between roughly 0.1 and 10 GeV^2. In particular, for reasonable values of the axial coupling one predicts that the asymptotic value of S_G is yet smaller than the presently reported one.

In conclusion, nonperturbative evolution of the Gottfried sum due to bound-state emission in an intermediate energy range $0.3\ \text{GeV}^2 \leq |Q| \leq 3$ GeV^2 is capable of reproducing the difference between the quark model value and the experimental result taken at $Q^2 = 4\ \text{GeV}^2$. The quark model value holds at a low scale, and is reduced by nonperturbative generation of a flavor-asymmetric sea, generated through a nonperturbative generalization of the Altarelli–Parisi equations. The required asymmetry follows from the low-energy symmetry structure of QCD, as reflected in the spectrum of pseudoscalar mesons, and is due to the anomalous breaking of flavor singlet chiral symmetry. In this sense, the data on the Gottfried sum are displaying an effect of the axial anomaly. Because the nonperturbative evolution flattens, the value of S_G thus obtained is preserved by further QCD evolution: thus, nonperturbative QCD effects originating in the intermediate-energy region leave a spur even in the asymptotic limit.

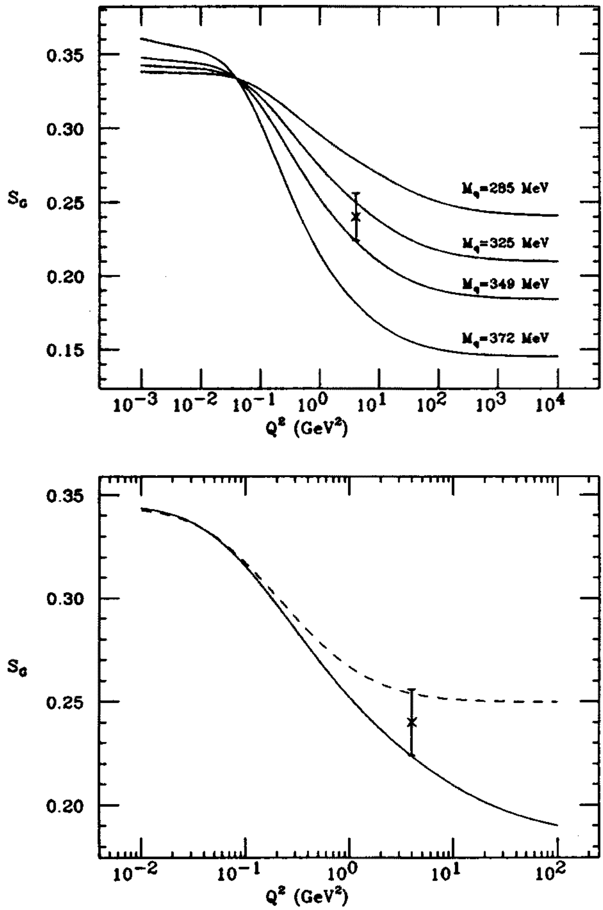


Fig. 7. Scale dependence of S_G computed from Eq. (31): a) variation with the constituent quark mass; the axial coupling is set to a maximal value; b) $M_q = 350$ MeV; full curve: maximal axial coupling; dashed curve: no axial coupling. The experimental point Eq. (11) is also shown.

3. The sigma term puzzle

We are finally left with assumption 3, Eq. (5). The current knowledge of the nucleon's structure functions is in good agreement with this expectation: for example, at a scale of Q^2=4 GeV^2 strange quarks carry about 2-3% of the nucleon's momentum [21] (and accordingly much less at the nucleon's scale, since singlet anomalous dimensions are generally large and in this case threshold effects should also be relevant). This suggests that the matrix element $\langle p|\bar{s}s|p\rangle$ should be roughly[2] of order 1 % of $\langle p|\bar{u}u|p\rangle$. However, there is an old piece of data which seems to contradict this conclusion; this is known as the puzzle of the pion-nucleon sigma term [7]. The pion-nucleon sigma term is defined as the nucleon matrix element of the light quark mass term in the QCD Hamiltonian:

$$\sigma \equiv \langle N|\bar{m}(\bar{u}u + \bar{d}d)|N\rangle \,, \tag{34}$$

where $\bar{m}$ is the average light quark mass, *i.e.*, $\bar{m} = (m_u + m_d)/2$. The value of this operator can be extracted from experiment in two different ways; it would seem that the only way the results of the two determination can be made consistent with each other is by violating assumption 3. We show now, however, that a more careful analysis [22] reveals that this is not necessarily the case.

The most direct way (at least theoretically) of determining the value of the sigma term is to observe that scalar quark densities appearing on the r.h.s. of Eq. (34) can be obtained by commuting a pseudoscalar charge with a pseudoscalar density according to the basic current algebra relation (23)

$$\left[\bar{\psi}\gamma^0\gamma_5\tau^a\psi(x), \bar{\psi}\gamma_5\tau^a\psi(y)\right] = -\bar{\psi}\psi(x)\delta^{(4)}(x-y) \,, \tag{35}$$

where ψ are quark isospinors and τ^a are Pauli isospin matrices. The pseudoscalar current is in turn related to the pion interpolating field ϕ_{π^a} by [23]

$$\partial_\mu\bar{\psi}\gamma^\mu\gamma_5\tau^a\psi(x) = f_{\pi^a} m^2_{\pi^a}\phi_{\pi^a} \,, \tag{36}$$

where a labels the pion according to its isospin and f_{π^a} and m_{π^a} are the pion's decay constant and mass. Using Eqs (35), (36) the matrix element on the r.h.s. of Eq. (34) can be related through straightforward current algebra manipulations [23, 24] to a pion-nucleon matrix element which, in turn, is related to the isospin-even pion-nucleon scattering amplitude at vanishing momentum transfer. This, up to several subtle details of the extraction of the matrix element from the data, provides a direct unambiguous determination of σ as defined by Eq. (34), with the result [25]

$$\sigma = 45 \pm 10 \ \text{MeV} \,. \tag{37}$$

On the other hand, if we assume Eq. (5) (*i.e.*assumption 3) to hold, we can equate the singlet and SU(3)-octet sigma terms:

$$\sigma \approx \sigma_8 \equiv \bar{m}\langle N|(\bar{u}u + \bar{d}d - 2\bar{s}s)|N\rangle \,. \tag{38}$$

[2] It can actually be shown rigorously [15] by defining parton distributions in terms of QCD light-cone operators that the forward matrix element of the operator $\bar{\psi}_i\psi_i$ is to leading perturbative order equal to the first moment of the distribution of quarks plus antiquarks of flavor i.

The reason why such an identification is interesting is that σ_8 is proportional to the matrix element of the octet portion $\mathcal{H}_8$ of the QCD Hamiltonian, which coincides with the octet portion of the mass term since SU(3) flavor symmetry is only broken by masses in QCD:

$$\sigma_8 = \frac{3}{1 - m_s/\bar{m}} \langle N|\mathcal{H}_8|N\rangle \tag{39}$$

$$\mathcal{H}_8 = \tfrac{1}{3}\left(\bar{m} - m_s\right)\left(\bar{u}u + \bar{d}d - 2\bar{s}s\right) . \tag{40}$$

Because in turn the matrix element of the octet part of the Hamiltonian provides SU(3) breaking to first order in perturbation theory, it can be measured from SU(3) mass splittings, for example in the baryon octet.

Indeed, it is well known that in the nonrelativistic quark model one assumes the mass splittings in the baryon octet to be given by first order in perturbation by an operator $\mathcal{O}_8$ which transforms as an octet under flavor SU(3) (identified with the octet component of the strong Hamiltonian). Group theory then leads to [24]

$$M_B = M_0 + M_1\langle B|Y|B\rangle + M_2\langle B|\left[I(I+1) - \tfrac{1}{4}Y^2\right]|B\rangle , \tag{41}$$

where I and Y, respectively, are isospin and hypercharge operators. Eq. (41) expresses the masses of the all the octet baryons in terms of the three free parameters M_i and is in excellent agreement with experiment; in particular the Gell-Mann–Okubo mass formula, which is a consequence of it, is one of the great successes of the quark model. The nucleon matrix element of the symmetry breaking operator $\mathcal{O}_8$ can also be determined from Eq. (41) using SU(3) algebra:

$$\langle N|\mathcal{O}_8|N\rangle = M_1 - \frac{M_2}{2} = M_\Lambda - M_\Xi = \frac{2M_N - M_\Xi - M_\Sigma}{3} . \tag{42}$$

Identifying thus the symmetry breaking operator $\mathcal{O}_8$ with the octet component of the QCD Hamiltonian $\mathcal{H}_8$, and using the known [26] value of the (current) quark mass ratio $m_s/\bar{m} = 25 \pm 5$, Eqs (39), (42) determine the octet sigma term as $\sigma_8 = 25 \pm 5$ MeV, in blatant contradiction with the value of σ Eq. (37), *i.e.* with the identification of σ and σ_8 Eq. (38). This discrepancy is known as the sigma term puzzle.

The most simple-minded conclusion, namely that σ_8 is much smaller than σ just because the nucleon matrix element of $\bar{s}s$ is large leads to rather dramatic (and hard to believe) consequences: because the strange quark is so much heavier than the light ones, it would follow that about one third of the nucleon's mass is carried by strange quarks, in contradiction to the parton model results [21] discussed above. Alternative explanations, based on a failure of first-order perturbation theory (in which case it's Eq. (41) and its consequences which do not hold) are extremely hard to reconcile with the success of the corresponding quark model formulas. It is interesting to observe that, once again, effective models of the nucleon have no difficulty in accommodating this state of affairs: in the Skyrme model [27] indeed assumption 3 is violated and σ_8 and σ differ because of a large strange contribution, whereas in the bag model [27] assumption 3 is valid but first-order perturbation theory fails.

Upon closer examination [22], however, the argument given above is seen to contain a loophole: even though it is a true empirical fact that octet mass splittings are accurately described by the first order perturbative formula Eq. (41), how do we know that the operator $\mathcal{O}_8$ which appears in this formula is the same as the operator $\mathcal{H}_8$ which is related to the sigma term by Eq. (39)? Of course, classically this is necessarily the case: mass splittings are given by the symmetry-breaking portion of the Hamiltonian, and that's what $\mathcal{H}_8$ is. However in a second-quantized field theory this needs not necessarily to be the case: quantum corrections may change the symmetry structure, in general, hence also the pattern of symmetry breaking (through a quantum anomaly). Now, it is well known that this is precisely what happens for masses in the strong interactions: the masses which enter quark model formulas are constituent quark masses, which for light quarks differ by two orders of magnitude from the current quark masses which appear in the fundamental Lagrangian (and Hamiltonian) of the theory. Hence, what $\mathcal{O}_8$ should be identified with is the octet piece of the *constituent* quark mass term, *i.e.*, with the octet piece of the effective low energy Hamiltonian of QCD [29], which is generally quite different from the fundamental QCD Hamiltonian. Thus we have the possibility of escaping the above dilemma if it just so happens that the octet portion of constituent quark masses and current quark masses are different: in such a case Eq. (38) still holds, along with first-order perturbation theory and its quark model consequences, but, at least in principle, octet splittings calculated from $\mathcal{O}_8$ could be smaller than those calculated from $\mathcal{H}_8$, which would explain the above discrepancy without invoking a violation of assumption 3.

This explanation is actually less contrived than it may seem at first. We can see this by noting that actually the relationship between the matrix elements of $\mathcal{O}_8$ and $\mathcal{H}_8$ is an exact consequence of the so-called conformal anomaly equation. This is an expression for the trace of the energy-momentum tensor $T^{\mu\nu}$ which is true at the operator level in QCD:

$$T^{\mu}{}_{\mu} = (1+\gamma_m)\sum_i m_i \bar{\psi}_i \psi_i + \frac{\beta(\alpha_s)}{4\alpha_s} G_a^{\mu\nu} G^a_{\mu\nu}\,, \tag{43}$$

where the sum runs over all quark flavors, $G_a^{\mu\nu}$ is the gluon field strength,

$\beta = \frac{d\alpha_s(\mu)}{d\ln\mu}$ is the beta-function for the strong coupling $\alpha_s = \frac{g^2}{4\pi}$, and $\gamma_m(\mu) = -\frac{d\ln m(\mu)}{d\ln\mu}$ is the mass anomalous dimension. The last term on the r.h.s. of Eq. (43) is due to the conformal anomaly; the anomalous dimension γ_m is present because the anomaly term and the mass term in Eq. (43) are not separately scale invariant (*i.e.*, they mix upon renormalization), while the energy momentum tensor is (up to surface terms). Classically, both γ_m and the gluonic contributions would be absent.

Taking the baryon matrix element of Eq. (43) and using Lorentz invariance [23] (*i.e.*, essentially the fact that the trace of the energy-momentum tensor is a measure of the violation of scale invariance) it follows that the baryon mass M_B is given by

$$\langle B|\,(1+\gamma_m)\sum_i m_i\bar{\psi}_i\psi_i + \frac{\beta(\alpha_s)}{4\alpha_s}G_a^{\mu\nu}G^a_{\mu\nu}|B\rangle = M_B\,. \tag{44}$$

This shows explicitly that indeed mass splittings are linear in the expectation value of a certain operator, as in lowest order perturbation theory: because of the conformal anomaly, the linear "approximation" is actually exact; quantum corrections appear as an additive contribution to the naive first order perturbative formula. However, the operator which ought to be used is that on the r.h.s. of Eq. (44). Otherwise stated, Eq. (44) shows that

$$\mathcal{O}_8 = \left[(1+\gamma_m)\sum_i m_i\bar{\psi}_i\psi_i + \frac{\beta(\alpha_s)}{4\alpha_s}G_a^{\mu\nu}G^a_{\mu\nu}\right]_{(8)}\,, \tag{45}$$

where the subscript (8) indicates the SU(3) octet component of the operator in square brackets.

It follows that the baryon mass M_B is in general given by the sum of a current contribution M_B^C and a dynamical contribution M_B^D:

$$M_B = M_B^C + M_B^D\,; \tag{46}$$

$$M_B^C = \langle B|\sum_i m_i\bar{\psi}_i\psi_i|B\rangle\,, \tag{47}$$

$$M_B^D = \langle B|\frac{\beta(\alpha_s)}{4\alpha_s}G_a^{\mu\nu}G^a_{\mu\nu} + \gamma_m\sum_i m_i\bar{\psi}_i\psi_i|B\rangle\,. \tag{48}$$

The octet portion of $\mathcal{H}_8$ provides the current mass splittings of M_B^C, while the octet portion of $\mathcal{O}_8$ provides the full splittings of M_B and it reduces to $\mathcal{H}_8$ only in the limit in which both γ_m and the gluonic term can be neglected. It should be stressed that this statement is an exact consequence of QCD.

Superficially, one may think that the gluonic operator in Eq. (45) should be flavor singlet and thus should not contribute to mass splittings. In this case, one would be left with γ_m, which is however not quite negligible, since it may be computed in perturbation theory with the result [26] $\gamma_m[Q = 1 \text{ GeV}] = 0.27$.[3] This, however, is not possible since, as already mentioned, the gluonic contribution is not separately renormalization-group invariant, thus, by mixing with the fermionic component it acquires a nonsinglet portion. Indeed, Eq. (46) is the unique renormalization-group invariant decomposition of the baryon mass.

In fact, we can give two separate qualitative reasons which suggest instead that indeed M_B^D should be anticorrelated to M_B^C. The first reason is based on the observation that a Ward identity argument implies that the *one-meson reducible* contributions to M_B^C and M_B^D must cancel:

$$(M_B^C)^{\text{omr}} + (M_B^D)^{\text{omr}} = 0\,. \tag{49}$$

Here M_B^{omr} are defined as the contribution to the matrix elements Eqs (46)–(48) where the operators which enter the definitions Eqs (47), (48) couple to the baryon state by first coupling to a meson which then couples to the baryon. Because such contributions are expected (by pole-dominance arguments) to provide the bulk of M_B^D, Eq. (49) suggests that M_B^D has a large component anticorrelated to M_B^C.

An independent reason is given by the observation that Eq. (44) can be viewed as a relation between current masses m_q^C (given by Eq. (47) assuming that the matrix element of $\bar\psi_i\psi_i$ is just the number of quarks plus antiquarks of flavor i) and constituent quark masses m_q (which add up to the baryon mass M_B); accordingly, the decomposition Eq. (46) determines the dynamical contribution to the constituent quark mass (*i.e.*, the difference between constituent and current mass) as a function of the current mass m_q^C:

$$m_q^D \equiv m_q(m_q^C) - m_q^C = m_q^D(m_q^C)\,. \tag{50}$$

A baryon mass splitting will coincide with the mass splitting computed at the current level, *i.e.*, neglecting the dynamical contribution M_B^D Eq. (48), only if m_q^D Eq. (50) turns out to be a flat function, *i.e.*, not to depend on the current mass, so that the current and constituent masses differ by a fixed, flavor independent amount. However, the qualitative behavior of the function $m_q^D(m_q^C)$ Eq. (50) is actually known. Indeed, when the current mass vanishes, the constituent mass reduces to roughly 300 MeV, *i.e.*, a third

[3] Notice that if this were the only effect of the conformal anomaly, the discrepancy discussed above would be even larger, since this extra term would make $\mathcal{O}_8$ larger than $\mathcal{H}_8$, rather than smaller as required to explain the puzzle.

of the mass of the nucleon, which is made of quasi-massless quarks; hence $m_q^D(m_q^C = 0) = 300$ MeV. When the current mass tends to infinity, quarks become quasi-static, and the current and constituent masses coincide; hence $\lim_{m_q^C\to\infty} m_q^D(m_q^C) = 0$. It follows that m_q^D is a *decreasing* function of the constituent mass, and not a flat function. This implies that constituent mass splittings are in general smaller than current mass splittings (so, for instance, the strange-light splitting is *smaller* for constituent masses than it is for current ones): the contribution of M_B^D to mass splittings is anticorrelated to that of M_B^D.

Both these qualitative arguments can actually be turned into quantitative (or at least semi-quantitative) ones [22]. Firstly, one can actually estimate $(M_B^D)^{\rm omr}$ by pole-dominance arguments, thus obtaining a rough estimate of the value of M_B^D. Also, it is possible to actually compute the function $m_q^D(m_q^C)$ by studying the full quark propagator: more specifically, it is possible to attempt an approximate resolution of the Schwinger-Dyson equation satisfied by the full quark propagator. This leads to a determination of the quark self-energy, thus also of the quark constituent mass; this in turn can be assumed to give the baryon mass by an additive relation (*i.e.*, the baryon mass is given by the sum of the masses of its constituent quarks). In either case, it is possible to determine ΔM, defined as

$$\Delta M \equiv \langle N|\mathcal{O}_8|N\rangle - \langle N|\mathcal{H}_8|N\rangle\,. \tag{51}$$

Even though the former determination is rather crude, and the latter is affected by an underlying theoretical uncertainty due to lack of knowledge of the QCD dynamics in the infrared (as reflected, for example, by the behavior of the strong coupling in the infrared limit), both determinations lead to a value $\Delta M \sim 150$ MeV (with large uncertainties, perhaps up to 50 %). Specifically, with [22] 30 MeV $\leq \Delta M \leq$ 250 MeV, and using Eq. (51) to relate Eq. (41) (mass splittings) to Eq. (39) (matrix elements of the sigma term) leads to the value 30 MeV $\leq \sigma_8 \leq$ 60 MeV, in perfect agreement with the value of σ Eq. (37), *i.e.* with the identification of σ and σ_8, consistently with assumption 3 (Eq. (5))

In conclusion, if one takes into account that even "gluonic" quantities such as the dynamical contribution to baryon masses can have a flavor nonsinglet component, it follows that mass splittings at the constituent and current level are not the same. This allows one to resolve the puzzle of the discrepancy between the two different determinations of the sigma term, without invoking a large strange content of the nucleon, and without violating the linear quark model mass formulas, such as those leading to the Gell-Mann–Okubo relation.

4. Conclusions

There are two main lessons to be learnt from this work, a phenomenological one and a theoretical one. The phenomenological lesson is that instances of discrepancy between experimental data and expectations based on the quark model are resolved once nonperturbative effects are taken into account. The theoretical lesson is that these effects appear both in the QCD evolution of operators, and in the computation of the operators themselves (in particular, in the determination of their symmetry structure). Such contributions, even though of course they originate in the gluon dynamics, depend strongly on the quantum numbers of the physical states under investigation and cannot be simply parametrized by universal flavor-singlet parameters, such as vacuum condensates. These two points, taken together, suggest that new data may actually lead us to a better understanding of the infrared dynamics of QCD and thus to go beyond the very successful but limited set of predictions which perturbative QCD allows to make.

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