The Spin Structure of the Baryons

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Abstract

Baryons are considered in the Nonrelativistic Quark Model (NQM) to be bound states of three valence quarks. Each quark has two possible spin eigenstates in the restframe of the baryon whose spin is fully carried by quarks. The baryon wavefunctions are connected through SU(6) symmetry rotations. For a long time, the measured magnetic moments of the baryons appeared to be in agreement with the NQM predictions.
However, recent experiments which are examining the spin structure of the baryons show the failure of several NQM predictions.
The so-called 'spin crisis' arose from the interpretation of the EMC deep inelastic scattering measurement of $\int g_1^p$ that the quark spins in the proton appeared to sum up to (almost) zero.
In this thesis it will be demonstrated that the spin problem is not a phenomenon restricted to quasi-massless current quarks in the high energy limit. Symmetry arguments are used to examine the baryon magnetic moments and reveal that we can observe massive but pointlike constituent quarks, with a characteristic mass ratio $m_u = m_d \simeq \frac{2}{3} m_s$. Surprisingly they do not contribute much to the baryon spin either. This analysis is free of the ambiguity arising from the $U_A(1)$ gluon anomaly which makes it impossible to calculate precisely the spin sum of the current quarks. One important finding in our analysis is the observation that the effects of SU(6) breaking hyperfine spin-spin interactions (which cause well-known splittings in the baryon masses) can be seen in the environment dependence of the constituent quark masses. The effective mass of a quark cannot be independent of its surrounding energy since the mass of the baryon is distributed amongst its constituents.
Consistent with the hypothesis that different quark masses do not impose SU(3) breaking on the baryon wavefunctions is the observation of induced 'second class' form factors. The way in which SU(3) breaking alters the $\frac{g_A}{g_V}$ ratios in semileptonic hyperon decays will be discussed and strong evidence for a new value of F/D is given, which is close to its SU(6) value. This value is derived independently from the baryon β-decays <u>and</u> from their magnetic moments. Dynamical models are discussed which might explain the observed polarised strangeness 'inside' the proton, and the almost vanishing quark spin sum.

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1. Introduction

1.1. Historical background

A truly elementary particle is pointlike and therefore, according to Dirac's relativistic quantum theory, has no anomalous magnetic moment. About sixty years ago the measurement of the anomalous magnetic moment of the proton was a first indication of the existence of a substructure of the nucleon. Just before this, the 'theoretical' discovery of isospin symmetry had suggested that the proton and the neutron are fundamental particles themselves.

Three decades later the quark model was proposed, after an increasing number of strongly interacting particles, with a clear pattern in their quantum numbers, had been detected. Only three different *flavours* with distinct quantum numbers (the *up*, *down* and *strange* quarks and antiquarks) were needed for the classification of all known hadrons (mesons and baryons). All baryons, including the nucleons, could apparently be interpreted as $|qqq>$ bound states of three quarks, whereas the mesons are $|\bar{q}q>$ quark-antiquark configurations.

The mathematical description of these configurations is given by the SU(3) Lie-group. The strong force binds quarks and antiquarks in an invariant manner under a global quark flavour transformation $\psi \rightarrow \psi' = e^{i\epsilon_a \lambda^a}\psi$, where $\psi = \begin{pmatrix} u \\ d \\ s \end{pmatrix}$ is a fundamental quark flavour triplet and λ^a are the generators of the SU(3) Lie-group. The

Gell-Mann matrices λ^a satisfy the SU(3) commutation relation $[\frac{\lambda^a}{2}, \frac{\lambda^b}{2}] = if^{abc}\frac{\lambda^c}{2}$. The SU(3) structure constants f^{abc} obey themselves the SU(3) commutation relation if they are defined as a set of matrices, and form the adjoint representation.

Quarks and antiquarks are spin-$\frac{1}{2}$ particles and the Lie-group which describes the spin transformation is SU(2). This is generated by the well-known Pauli matrices σ_i which satisfy the SU(2) commutation relation $[\sigma_i, \sigma_j] = i\epsilon_{ijk}\sigma_k$, where ϵ_{ijk} is the Levi-Civita symbol.

If the strong force is also invariant under the SU(2) spin transformation of quark- and antiquark-spins, all 35 generators σ, $\frac{\lambda}{2}$ and $\sigma \cdot \frac{\lambda}{2}$ generate the SU(6) algebra. In this case irreducible representations of this large group can be expected to manifest themselves as degenerate particle multiplets.

The Pauli principle underlies the classification of the hadronic states in terms of the elementary quarks and antiquarks. The Nonrelativistic Quark Model (NQM) [1] classifies all known hadrons by assigning additive quantum numbers to these states in the rest frame of the particles. In this frame the four-component Dirac-spinor can be replaced by the two-component Pauli-spinor, the quark spin being polarised parallel (or anti-parallel), i.e. $|q^{\uparrow(\downarrow)}>$, to the hadron spin $|h^{\uparrow}>$. The fundamental representation 6 of the SU(6) group in the NQM is given by $\psi^T = (u^{\uparrow}, u^{\downarrow}, d^{\uparrow}, d^{\downarrow}, s^{\uparrow}, s^{\downarrow})$. Having three quarks inside a baryon the decomposition $6 \otimes 6 \otimes 6 = 20 \oplus 56 \oplus 70 \oplus 70$ tells us that the 56-plet is flavour-spin symmetric. It contains the observed $J^P = \frac{3}{2}^+$ flavour decuplet and the $J^P = \frac{1}{2}^+$ flavour octet states; $56 = (10 \otimes 4) \oplus (8 \otimes 2)$. These lowest lying baryons in the mass spectrum indicate that there is no significant difference between two ways in which SU(6) is broken: The magnitude of the spin induced interaction which results in the mass difference between the flavour decuplet and the flavour octet states is of the same order as the mass splitting within one of these irreducible representations of SU(3) due to unequal quark masses. Being subject to an attractive strong force, the quarks in their restframe of a groundstate baryon are localised in an s-wave with no angular momentum $L^P = 0^+$. The spin-statistics problem of having a flavour-spin symmetric wavefunction of the 56-plet and a symmetric spacefunction was solved with the discovery of the colour-degree of freedom. Being totally antisymmetric in the wavefunction quarks have a hidden three-valued colour quantum number $N_{colour} = 3$. In the NQM the baryon wavefunction is sufficiently described by its spin-flavour combination. In the case of the proton we have:

$$\begin{aligned} |p^\uparrow> = \frac{1}{\sqrt{18}}(2|u^\uparrow u^\uparrow d^\downarrow > &+2|u^\uparrow d^\downarrow u^\uparrow > +2|d^\downarrow u^\uparrow u^\uparrow > \\ &- |u^\uparrow u^\downarrow d^\uparrow > -|u^\downarrow u^\uparrow d^\uparrow > -|u^\uparrow d^\uparrow u^\downarrow > \\ &- |u^\downarrow d^\uparrow u^\uparrow > -|d^\uparrow u^\uparrow u^\downarrow > -|d^\uparrow u^\downarrow u^\uparrow >) \end{aligned} \tag{1.1}$$

Clearly the spin of the groundstate baryons is 100% the sum of all quark spins. One of the great successes of the NQM is the parameter free prediction of the ratio of the magnetic moments between the proton and the neutron:

$$\frac{\mu_p}{\mu_n} = -\frac{3}{2} \tag{1.2}$$

This result is in acceptable agreement with the data

$$\frac{\mu_p}{\mu_n}|_{exp} = -\frac{2.79}{1.91} \tag{1.3}$$

If the quarks in the NQM are bare Dirac particles with no anomalous magnetic moments ($\kappa_q = 0$) then the absolute values of the magnetic moments μ_p, μ_n and

μ_Λ of the proton and the Λ set the scale for the up-, down- and strange-quark masses m_u, m_d, m_s;

$$m_u = \frac{m_{proton}}{2.79} \simeq 340 MeV, \qquad \text{and (1.2)} \quad \Rightarrow \quad m_d \simeq 340 MeV$$

$$\frac{\mu_\Lambda}{\mu_p} = -\frac{1}{3}\frac{m_u}{m_s} = -0.22 \quad \Rightarrow \quad m_s \simeq 500 MeV \tag{1.4}$$

These values for the effective quark masses are attractive for estimating the baryon masses. However, with the constituent quarks having such a comparatively small mass, it is not clear how the non-relativistic nature of the NQM can be justified. Being confined by the sphere of a baryon to a region $|r| < R$, the magnetic moment of a quark with wavefunction ψ and electrical charge e_q is

$$\mu_q = \frac{1}{2}\int_{|r|<R} d^3r r \times (\bar{\psi}\gamma\psi)e_q \tag{1.5}$$

and the anomalous magnetic moment of a quark is not expected to vanish, unless

$$m_q >> \frac{1}{R} \tag{1.6}$$

with $\frac{1}{R}$ being the characteristic momentum of the quark confined to R. In this context, however, it is worth noting that the whole problem of the quark confinement is still to be solved.

After the introduction of the colour degree of freedom, the confining strong forces between the coloured quarks had also to be colour dependent, because only colour singlets are observable.

About twenty years ago the theory for the dynamics of the strong interaction, namely Quantum Chromo Dynamics (QCD), was formulated. Like the electroweak theory, QCD is also a renormalizable local gauge theory. Its gauge group is $SU(3)_{colour}$. Its Lagrangian is usually written as:

$$\mathcal{L}_{QCD} = -\frac{1}{2} tr G_{\mu\nu} G^{\mu\nu} + \sum_{k}^{n_f} \bar{q}_k (i\gamma^\mu \mathcal{D}_\mu - m_k) q_k$$

where

$$\begin{aligned} G_{\mu\nu} &= \partial_\mu A_\nu - \partial_\nu A_\mu - ig[A_\mu, A_\nu] \\ \mathcal{D}_\mu &= (\partial_\mu - igA_\mu) \\ A_\mu &= \sum_{a=1}^{8} A_\mu^a \frac{\lambda^a}{2} \quad \text{with} \quad [\frac{\lambda_a}{2}, \frac{\lambda_b}{2}] = i f^{abc} \frac{\lambda_c}{2} \end{aligned} \tag{1.7}$$

where λ^a are again the Gell-Mann matrices that satisfy the SU(3) commutation relation. Various experiments have shown the existence of the colour degree of freedom of the strong interaction. The decay rate of $\pi^0 \rightarrow 2\gamma$ is proportional to N_c, and the cross-section $\sigma(e^+e^- \rightarrow hadrons)$ is proportional to N_c^2. These experiments confirm $N_c = 3$. Also the anomaly condition of the Standard Model that the sum of all electric charges from all quarks and leptons in one family vanish requires $N_c = 3$.

The effective coupling of the strong interaction is governed by the Renormalisation Group (RG) β-function. The running coupling constant α_s of the strong interaction is well defined in terms of a perturbative formula, but its validity is restricted to the ultraviolet region where all particles are far away from their mass shell. The strong coupling is given by

$$\alpha_s(Q^2) = \frac{4\pi}{(11 - \frac{2}{3}n_f) ln \frac{Q^2}{\Lambda^2}} \tag{1.8}$$

where Λ_{QCD} is the fundamental momentum scale of the theory and n_f is the number of flavours. The asymptotic freedom of the strong interaction means that α_s decreases with increasing Q^2, and it is plausible that massless quarks and gluons are confined. Λ_{QCD} is not yet precisely determined, but it is measured to be in the range of 200 MeV. With respect to the light flavours, the *up*, *down* and *strange* quarks, the QCD Lagrangian $\mathcal{L}_{QCD}$ displays a chiral symmetry: at high energies the masses of these quarks, being smaller than Λ_{QCD}, are negligible. Due to the missing rest frame of massless particles, no transformation, which is to say no interaction, is possible in which the helicity of the particle is changed. These massless quarks are obviously not the above-mentioned massive quarks which are observable via the magnetic moments in equation (1.4). Instead of the massive constituent quarks with $m_q > \Lambda_{QCD}$ one speaks of quasi massless current quarks with $m_q < \Lambda_{QCD}$, since conserved quark currents can be observed.

With q_L and q_R being the left- and right-handed quark-fields

$$q_L = \frac{1}{2}(1-\gamma_5)q, \qquad q_R = \frac{1}{2}(1+\gamma_5)q$$

the QCD-Lagrangian

$$\mathcal{L}_{QCD} = \mathcal{L}_0 + \mathcal{L}_1 \tag{1.9}$$

contains one chiral symmetric part

$$\mathcal{L}_0 = i\bar{q}_L \mathcal{D}_\mu \gamma^\mu q_L + i\bar{q}_R \mathcal{D}_\mu \gamma^\mu q_R$$

and the chiral symmetry breaking part

$$\mathcal{L}_1 = m_u(\bar{u}_L u_R + \bar{u}_R u_L) + m_d(\bar{d}_L d_R + \bar{d}_R d_L) + m_s(\bar{s}_L s_R + \bar{s}_R s_L)$$

Noether's theorem gives the connection between symmetries and conservation laws; namely that any continous symmetry transformation which leaves the Lagrangian invariant implies the existence of a conserved current. In the case that $m_u = m_d = m_s$ there are eight conserved vector currents

$$V^a_\mu(x) = \bar{q}(x)\gamma_\mu(\frac{\lambda^a}{2})q(x) \ , \qquad \partial^\mu V^a_\mu = 0 \tag{1.10}$$

and in the case that $\mathcal{L}_1 = 0$ (i.e. $m_u = m_d = m_s = 0$) there are eight additional conserved axial vector currents

$$A^a_\mu(x) = \bar{q}(x)\gamma_\mu\gamma_5(\frac{\lambda^a}{2})q(x) \ , \qquad \partial^\mu A^a_\mu = 0 \tag{1.11}$$

The additional U(1) vector symmetry displayed by $\mathcal{L}_o + \mathcal{L}_1$ corresponds to the conserved baryon number current. The U(1) axial vector symmetry, however, does not yield a conserved current even if $m_q = 0$. The flavour singlet axial vector current

$$A^0_\mu(x) = \bar{q}(x)\gamma_\mu\gamma_5 q(x) \tag{1.12}$$

is not conserved due to the gluon anomaly:

$$\partial^\mu A^0_\mu(x) = \frac{g^2}{4\pi} trGG^* \tag{1.13}$$

where $G^{\mu\nu}$ is the gluon tensor matrix and $G^*_{\mu\nu} = \epsilon_{\mu\nu\rho\sigma}G^{\rho\sigma}$ is its dual.

The chiral SU(3) $\otimes$ SU(3) quark symmetry, however, is not realized at the macroscopic level of the physical groundstate. The existence of eight pseudoscalar mesons, all of which are comparatively light, suggests the Spontaneous Symmetry Breakdown (SSB) of the chiral symmetry. The flavour octet chiral charges

$$Q^{5a} = \int d^3x A_0^a(x) \tag{1.14}$$

do not annihilate the vacuum, in fact the Nambu-Goldstone (NG)-bosons $\pi^a(x)$ have direct couplings to the broken axial charges Q^{5a} and currents A_μ^a:

$$\begin{aligned} &< 0|A_\mu^a(0)|\pi^b(p) > = \delta^{ab} f_\pi p_\mu \\ \text{and} \quad &< 0|\partial^\mu A_\mu^a(0)|\pi^b(p) > = f_\pi m_\pi^2 < 0|\phi^a(0)|\pi^b(p) > \end{aligned} \tag{1.15}$$

This leads us to the PCAC (Partially Conserved Axialvector Current) hypothesis

$$\partial^\mu A_\mu^a = f_\pi m_\pi^2 \phi^a \tag{1.16}$$

which connects the weak currents A_μ^a and the strong interacting pion fields ϕ^a in an operator relation. The $J^P = 0^-$ spinless π's, K's and the η particles may be regarded as the NG-bosons $\pi^a(x)$, and are the necessary degrees of freedom left over from the SSB. Since the SU(3) vacuum is not spontaneously broken, the vacuum is taken to be SU(3) symmetric, and consequently the decay constants are equal:

$$f_\pi = f_K = f_\eta =: f \tag{1.17}$$

This explanation for the existence of the pseudoscalar mesons is essentially different from the one given by the NQM. In the NQM these particles are taken to be the 'astonishingly light' $J^P = 0^-$ bound states of constituent quarks, which have to be compared with the much heavier $J^P = 1^-$ vector-mesons.

A longstanding question has been, how to calculate the properties of the hadrons from first principles. The 'success' of the NQM cannot be sufficiently understood in the more general terms of QCD. But this no longer seems to be a weakness of QCD. At present, particle data tend to suggest that it is not the NQM which has to be derived from QCD. Instead, the NQM must be justified as being the correct model for the description of hadronic properties. It is already part of the popular literature [2] that the lowest lying groundstate of the baryons, namely the proton, is insufficiently understood in the framework of the NQM. The target of many high precision measurements in recent years, the structure of the proton, has been proven to be quite complicated.

Of course, as we have noted beforehand, the quarks in the NQM are not the quarks in $\mathcal{L}_{QCD}$. Even today it is not clear what exactly the constituent quarks are and how they are generated from the current quarks. Generally, the massive constituent quarks are considered to be the product of the SSB of the chiral symmetry. We have this quite vague picture in mind when we talk of constituent quarks. However, constituent quarks are physical entities which are observable at a scale below the scale where the chiral symmetry breaking takes place. One can expect to observe them indirectly via the measurements of the baryon masses or magnetic moments. In the context of the discussion of the baryon magnetic moments a definition of constituent quarks via their masses (1.4) is suggested in the framework of the NQM. A more general discussion will take place in chapter 4. There we do not refer to an absolute value of the constituent quark mass but to their flavour mass ratio $m_u : m_d : m_s$.

A model which describes the dynamical generation of a constituent quark mass and another model that takes constituent quarks as fundamental degrees of freedom in an effective field theory will be discussed in chapter 5. In a field theory the absolute number of quarks plus antiquarks is not conserved. The conservation of the baryon number however, implies that the difference between quarks and antiquarks is conserved. Therefore, in the case that we want to emphasize the flavour content of a baryon we talk of valence quarks q_V, $u_V := u - \bar{u}, d_V := d - \bar{d}, s_V := s - \bar{s}$.

Many discrepancies between the predictions of the NQM and experimental results are related to the problem of the spin structure of the baryons.

SU(6) contains the SU(2) subgroup (which may be generated by the Pauli matrices σ_j) and one can think of this as the 'intrinsic spin' of the hadron, or the quark system's total spin, before coupling with the orbital angular momentum in the system. That this is not a good decay symmetry within the NQM has been observed by looking at the NQM forbidden decays like

$$\begin{aligned} &\Delta_{(J^P=\frac{3}{2}^+)} \to N_{(J^P=\frac{1}{2}^+)} \otimes \pi_{(J^P=0^-)} \\ \text{and} \quad &\rho_{(J^P=1^-)} \to \pi_{(J^P=0^-)} \otimes \pi_{(J^P=0^-)} \end{aligned} \tag{1.18}$$

But in a relativistical extension of the NQM the intrinsic spin of the hadrons is conserved in these decays. The trick is to Lorentz-boost the quark q along the z-axis from rest to a frame with velocity $|v_z| = \tanh \omega$

$$q \to q' = e^{\frac{1}{2}\omega\alpha_3} \tag{1.19}$$

where $\alpha_i = \begin{pmatrix} 0 & \sigma_i \\ \sigma_i & 0 \end{pmatrix}$. This boost preserves the up-down spin quantization along the z-axis, but because α_3 does not commute with the other two generators of SU(2) transitions between different representations of that group can occur. The spin group which commutes with α_3 is found to be $SU(2)_W$ with the generators $(1, \beta\sigma_x, \beta\sigma_y, \sigma_z)$, where $\beta = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}$ is an intrinsic parity operator. Quarks q and antiquarks $\bar{q}$ transform differently under β

$$\beta q = +q \ , \qquad \beta \bar{q} = -\bar{q} \tag{1.20}$$

That is the reason why baryon states made of qqq transform in the same way for $SU(6)_W$ as they do for SU(6), but mesons being $q\bar{q}$ states transform differently. Under the assumption that π is a bound state made of constituent quarks, the above decay processes are allowed under $SU(6)_W$:

$$\begin{aligned} &\Delta_{(W=\frac{3}{2})} \rightarrow N_{(W=\frac{1}{2})} \otimes \pi_{(W=1)} \\ \text{and} \quad &\rho_{(W=0)} \rightarrow \pi_{(W=1)} \otimes \pi_{(W=1)} \end{aligned} \tag{1.21}$$

Another relativistic correction to the NQM is normally applied when explaining the measured 25% reduction from the NQM value $\frac{5}{3}$ of the axial-to-vector coupling constant ratio $\frac{g_A}{g_V}|_{n\rightarrow p} \simeq \frac{5}{4}$ of the neutron β-decay. In the NQM the quarks are at rest in the hadron rest frame and the four-component Dirac quark spinor $q(x)$ approximates to a two-component Pauli spinor χ_i. Because the quarks are believed to be confined to a sphere of radius R, their internal motion due to the uncertainty relation cannot be neglected. In fact, their momentum is expected to be of the order of their effective constituent mass:

$$(< p_i^2 >)^{1/2} \sim \frac{1}{R} \simeq 300MeV \tag{1.22}$$

The Dirac-spinor is written as

$$q_i(x) = \begin{pmatrix} \chi_i(x) \\ \frac{\sigma_i \cdot p_i}{2m_q} \chi_i(x) \end{pmatrix} \tag{1.23}$$

and the corresponding reduction of $\frac{g_A}{g_V}|_{n\rightarrow p} = \frac{5}{3}$ is generated by boosting $p_z \rightarrow \infty$, and is proportional (equation 6.97 in Close [1]) to

$$\sum_i \left(1 - \frac{p_i^2}{4m_q}\right) \tag{1.24}$$

(It is however not clear why the effective quark mass m_q should remain Lorentz-invariant.) A slightly different explanation for this reduction is given in the framework of the Melosh transformation for a free quark model. The Melosh transformation is similar in its form to the above picture because it requires the same combination of quark spin-flip and change in the quark's orbital angular momentum. The operator V for this transformation acts on a single quark

$$V \propto Z = i(\sigma_+ L_- - \sigma_- L_+) \tag{1.25}$$

where $\sigma_\pm, L_\pm$ are the spin and orbital angular momentum raising and lowering operators. The identity of the Melosh transformation with the Foldy-Wouthuysen transformation in an interaction free theory has been proven by Bell [3], who also pointed out in the same article, that the weakness of SU(6) lies in its reference frame dependence. For example, the magnetic moments of the baryons can only be measured in a frame which is not at rest, due to a non-vanishing momentum transfer. But why are they seemingly best explained in the NQM?

There is another related question. Having explained the reduction of $\frac{g_A}{g_V}|_{n\to p}$ from its NQM value (albeit with the above reservations) one still has to explain why the $\frac{g_A}{g_V}|_{\Sigma^-\to n}$ value of the β-decay of the Σ^- is not 25% smaller than its NQM value $-\frac{1}{3}$. In general, the fit of the corresponding values for the β-decays of the hyperons, related through the algebra of the axial vector currents, turns out to be relatively poor (see chapter 3). The resulting F/D ratio, which is a characteristic for the axial SU(3) symmetry, is quite far below its NQM value. This possibly arises from a misunderstanding of the spin structure of the baryons. Different schemes have been proposed in order to resolve these discrepancies, but no general consensus, in favour of one specific explanation, has emerged.

This is true also with respect to the big surprise concerning the spin structure of the proton, coming from the EMC experiment.

The experimental result of the 'quark spin sum measurement' by the European Muon Collaboration (EMC), was published about four years ago [4], and the interpretation of the EMC data suggests that strange quarks, which do not even occur in the NQM inside a nucleon, carry a considerable and a negative (!) portion of the proton spin. In addition, the astonishing observation that the quark spin sum is compatible with zero and definitely not close to one (the NQM value) led to the so-called 'spin crisis' [5].

$$a_0 s_\mu =< p|\bar{q}(x)\gamma_\mu\gamma_5 q(x)|p >\simeq 0 \tag{1.26}$$

where s_μ is the polarisation four-vector. Are such discrepancies merely phenomena related to the properties of current quarks, the basic quanta of QCD? This would imply that a solution of the spin crisis simply requires a proper understanding of the gluon anomaly (1.13). In this case, the problem of the proton spin would be solved by understanding the constituent quark spin structure [6].

Besides the data of the EMC experiment the recent high precision measurements of the magnetic moments of the baryons [7] (which are expected to be proportional to the constituent quark spin expectation values) display up to 37% deviations from the NQM values. This challenges the basis upon which the NQM rests.

In addition to all of the above there are the measurements of unexpected asymmetries in the polarized proton antiproton scattering [8], the unexpected polarized inclusive $\bar{\Xi}^+$ production [9] and the similarly unexpected difference in the electromagnetic polarizabilities of proton and neutron [10]. These all illustrate the need to overhaul the NQM. Experimental hints, that the corrections to the NQM may arise from (polarized) strangeness in the proton, have been given (by neutrino-proton [11]) and by pion-proton scattering [12].

1.2. Focus and organisation of the thesis

It is most challenging for a D.Phil. student to embark on a research project which embraces outstanding actuality, immense popularity and rather great generality. But the prospect of becoming a true specialist in a specific field by working on such a general topic like the spin structure of baryons is unfortunately less likely. The spin structure of the proton is of interest for the high-energy particle physicist as well as for the low-energy hadron and nuclear physicist, for the theoretician interested in the subtle formalism and the very mathematical aspects of anomalies and topologies, as well as for the phenomenologist interested in the physical explanation of particle properties. Because it is not yet known which 'part' of physics will be most decisive for giving the ultimate explanation of this problem, all kinds of models and ideas seem initially to be of equal importance.

However today, about four years after the EMC published its result [4], several important achievements in the general understanding of the problem allow us to focus on more specific problems. The first step was to recognise the existence of the gluon anomaly in the QCD-improved parton picture and the resulting modification of the quark spin sum [13,14][15]. Nevertheless, there are several reasons why a purely perturbative explanation is not sufficient to explain large differences between current and constituent quarks. Altarelli points to the lack of any dynamical reason

for this discrepancy as the *real* problem [16].
The second step should be the calculation of the matrix element a_0 of the axial vector current of the flavour singlet via a 'generalised' Goldberger-Treiman (GT) relation for the U(1) channel. The recent and comprehensive analysis of Shore and Veneziano [17] has shown that any such calculation is theoretically not justifiable, unless specific, model dependent assumptions are made.

A change of focus is therefore the logical consequence from these first steps which will be summarised in the next chapter. The *real* problem should now be approached via the constituent quarks. Constituent quarks are perceived as the fundamental entities necessary to explain with their additive quantum numbers and masses the basic hadronic properties. In a fundamental article, based on very general arguments, Weinberg offered an explanation for a constituent quark being a pointlike Dirac particle with no anomalous magnetic moment [18]. Weinberg's argument actually deepened the *real* problem of Altarelli.
Because on the other hand, the data of the baryon magnetic moments could not be explained so far within their experimental accuracy, it was necessary to pursue the possibility that so far undiscovered information about the spin structure of the baryons is the reason for this discrepancy. This question has motivated my work about the magnetic moments, and my contribution in this context has been the discovery of a wrong assumption which underlies all previous analysis.
The central conclusion is that the constituent quark mass should not be considered to be environment independent. If the constituent quark mass is allowed to fluctuate within a few MeV for each baryon state, violation of isospin is unnecessary to extract valuable and gluon-anomaly free information on the spin structure of all the flavour octet baryons. General symmetry considerations allow a simple analysis of the magnetic moments and this will be the core of chapter 4.

The reliability of the symmetry arguments used in the context of the magnetic moments is based on our understanding of the anomaly-free flavour octet axial vector currents A^a_μ, which are of crucial importance for the link between high-energy and low-energy phenomena. Because $dim[\partial^\mu A^a_\mu] < 4$, these currents are protected from large scale dependent QCD corrections. Their systematic evaluation can give us an insight into the pattern of SU(3) flavour symmetry breaking. This pattern should be matched at current and constituent quark level. If the SU(3) breaking in the quark masses is not to imply the complete loss of predictive power (i.e. the baryons remain SU(3) eigenstates) consistency requires that the characteristic F/D value has to be the same at current and constituent quark level.
It has been argued that the evaluation of the EMC-experiment and especially the size of the s-quark polarisation depends crucially on the F/D value [19,20] and on our knowledge about SU(3) [21]. Because of this there follows a discussion of the symmetry breaking pattern of SU(3) in chapter 3. Arguments are presented for a new value of F/D, who's determination from the semileptonic baryon decays is consistent with the analysis of the baryon magnetic moments, and holds a key role in scaling the spin problem from the current to the constituent quark level.

SU(3) symmetry, which is broken at first order only by a different strange quark mass, cannot however explain a negative strange quark polarisation inside the nucleon. Another change of focus is required. The *real* problem is now to arrive at an (SU(6)-symmetric?) flavour-spin and space symmetric baryon wavefunction built out of constituent quarks under the constraint of an (almost?) vanishing quark spin sum. In chapter 5, some dynamical models are introduced, which will hopefully show that this is well defined aim of future research.

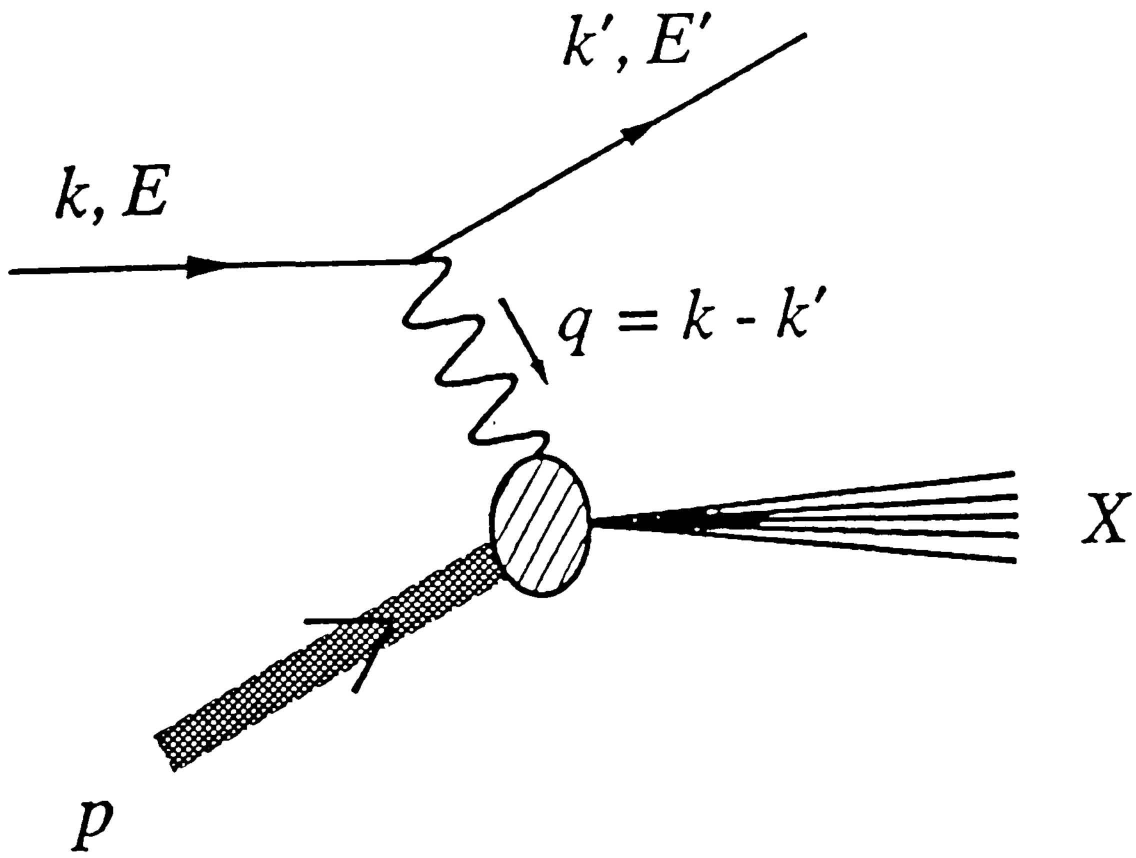
In chapter 6, the main aspects of the thesis are summarised. Strong arguments result from the analysis of the previous chapters and suggest another change of focus in the discussion of the 'spin problem'. The *real* problem defined by Altarelli was the question of how to explain dynamically a negative polarised strange current quark sea and a (probably) large and positive gluon polarisation inside the proton. Through the analysis of the axial vector current matrix elements and of the magnetic moments, it became obvious first that the spin problem is not restricted to the proton alone, and second that even at the constituent quark level the quark spin sum is close to zero. With this change of focus it becomes clear that the underlying dynamical explanation of the spin problem can be expected from a Lagrangian which has the constituent quarks as fundamental degrees of freedom.

The main arguments from chapter 3 and chapter 4 have been published in *Physics Letters B*[22,23] .

2. The Spin Dependent Structure Function g_1^p

2.1. The gluon contribution

In a classical deep inelastic scattering experiment (DIS), an incoming beam of leptons with energy E scatters off a fixed hadronic target. The high energy lepton interacts with the hadron target through the exchange of a virtual photon; the target hadron absorbs the virtual photon to produce the final state X. Only the energy and the direction of the scattered lepton are measured in the detector.



The kinematic variables necessary for the description of the process are

E The energy of the incident lepton.

k The momentum of the initial lepton.

E' The energy of the scattered lepton.

k' The momentum of the scattered lepton.

Ω The solid angle into which the outgoing lepton is scattered.

P The momentum of the target, $p = (M, 0, 0, 0)$, for a fixed target experiment.

M The target mass

$q = k - k'$ The momentum transfer in the scattering process.

$Q^2 = -q^2$

$\nu = E - E'$ The energy loss of the lepton.

$x = \frac{Q^2}{2M\nu}$ The Bjorken variable.

In leading electromagnetic order the differential cross section for inclusive charged lepton-nucleon scattering has the following form

$$d\sigma = \frac{\alpha_{em}^2}{\pi} \frac{1}{q^4} L_{\mu\nu} W^{\mu\nu} \frac{d^3k'}{(p \cdot k)E'} \tag{2.1}$$

Because the lepton tensor has a pointlike interaction with the virtual photon, the measurement gives information about the hadronic tensor $W_{\mu\nu}$ which cannot be computed directly from QCD, due to the non-perturbative effects in the strong interaction. However the parity and time reversal invariance of QCD and the current conservation ($q_\mu W^{\mu\nu} = 0$) allow us to decompose $W^{\mu\nu}$ into structure functions.

The most general form of $W_{\mu\nu}$, consistent with gauge invariance, is

$$\frac{1}{\pi}W_{\mu\nu}(q,p,s) = [-g_{\mu\nu} - \frac{q_\mu q_\nu}{q^2}]W_1(\nu,q^2) + \frac{1}{M^2}[p_\mu - \frac{\nu}{q^2}q_\mu][p_\nu - \frac{\nu}{q^2}q_\nu]W_2(\nu,q^2) + \frac{i}{M}\epsilon_{\mu\nu\rho\sigma}q_\rho\{s_\sigma[G_1(\nu q^2) + \frac{\nu}{M^2}G_2(\nu,q^2)] - (s\cdot q)p_\sigma\frac{1}{M^2}G_2(\nu,q^2)\} \tag{2.2}$$

For the conservation of the lepton current the terms proportional q_μ and q_ν may be omitted when contracting with $l_{\mu\nu}$. The most concise form for the hadronic tensor of a spin-$\frac{1}{2}$ target, with s^σ being the nucleon spin, is therefore:

$$W_{\mu\nu} = -F_1 g_{\mu\nu} + \frac{F_2}{p\cdot q}p_\mu p_\nu + \frac{ig_1}{\nu}\epsilon_{\mu\nu\lambda\sigma}q^\lambda s^\sigma + \frac{ig_2}{\nu^2}\epsilon_{\mu\nu\lambda\sigma}q^\lambda(p\cdot qs^\sigma - s\cdot qp^\sigma) \tag{2.3}$$

In the Bjorken limit $\nu, Q^2 \to \infty$ for $x = \frac{Q^2}{2M\nu}$ constant (M being the nucleon mass) the structure functions $F_i(x,Q^2)$ and $g_i(x,Q^2)$ display a scaling behaviour in QCD up to logarithmic corrections in Q^2. The dominant dependence of the structure functions on the dimensionless Bjorken variable x, is understood to be the scattering off pointlike particles inside the nucleon, called partons. In the high proton momentum limit $P \to \infty$, xP is seen as the momentum carried by the parton. The total cross-section is the convolution of the probability of finding a parton inside a nucleon times the photon-parton cross-section over all values of x (i.e. $0 \leq x \leq 1$). QCD predicts the evolution of the parton, i.e. the (anti)quark and the gluon distribution functions ($q(x,Q^2)$ and $G(x,Q^2)$ respectively) with Q^2. The Altarelli-Parisi (AP) equations [24,25] describe the Q^2 evolution of the distribution functions in terms of splitting functions $P_{ab}(x/x')$ which measure the probability for finding parton a with momentum fraction x 'inside' (i.e. radiated from) parton b and momentum fraction x'. Explicitely, the AP equations are for the unpolarised distributions

$$
\begin{aligned}
\frac{dq_i(x,t)}{dt} &= \frac{\alpha_s(Q^2)}{2\pi}\int_x^1 \frac{dx'}{x}[\sum_i q_i(x',t)P_{qq}(x/x') + G(x',t)P_{qG}(x/x')] \\
\frac{dG_i(x,t)}{dt} &= \frac{\alpha_s(Q^2)}{2\pi}\int_x^1 \frac{dx'}{x}[\sum_i q_i(x',t)P_{Gq}(x/x') + G(x',t)P_{GG}(x/x')]
\end{aligned}
\tag{2.4}
$$

where the sum is taken over all quarks and antiquarks of flavour i and $t = ln\frac{Q^2}{\mu^2}$. (μ^2 is a subtraction scale unspecified in perturbative QCD, arising from the RG)

To leading order the structure functions for inclusive unpolarised charged nucleon scattering are:

$$
\begin{aligned}
F_1(x,Q^2) &= \frac{1}{2}\sum_i e_i^2[q_i(x,Q^2) + \bar{q}_i(x,Q^2)] \\
F_2(x,Q^2) &= 2xF_1(x,Q^2)(1 + R(x,Q^2))
\end{aligned}
\tag{2.5}
$$

$R = \frac{\sigma_0}{\sigma_T}$ stands for the ratio between the longitudinal and transverse cross-sections and vanishes in the QCD uncorrected parton picture because of the spin-$\frac{1}{2}$ nature of the quarks. It is very difficult to determine the small value of R precisely and a value based on a QCD calculation is normally assumed in order to compare different data on F_2 from different experimental groups [26].

In the polarised case, only $\int g_1^p(x)dx$ has been evaluated by measuring the asymmetry A defined by

$$
A = \frac{d\sigma^{\uparrow\downarrow} - d\sigma^{\uparrow\uparrow}}{d\sigma^{\uparrow\downarrow} + d\sigma^{\uparrow\uparrow}} \tag{2.6}
$$

where the difference in the numerator is between the cross-sections from left-handed muons on a proton at rest, with its spin along the direction of the μ-beam and opposed to it.

$g_2(x)$ will not be measured until scattering with transversely polarised nucleons is performed. Jaffe pointed out [27,28] that it is not yet clear if the Burkhardt-Cottingham sum-rule [29]

$$\int_0^1 dx g_2(x, Q^2) = 0 \tag{2.7}$$

has to be fullfilled on general grounds. Shuryak and Vainshteyn [30] realised that in general in QCD, g_2 can be written as

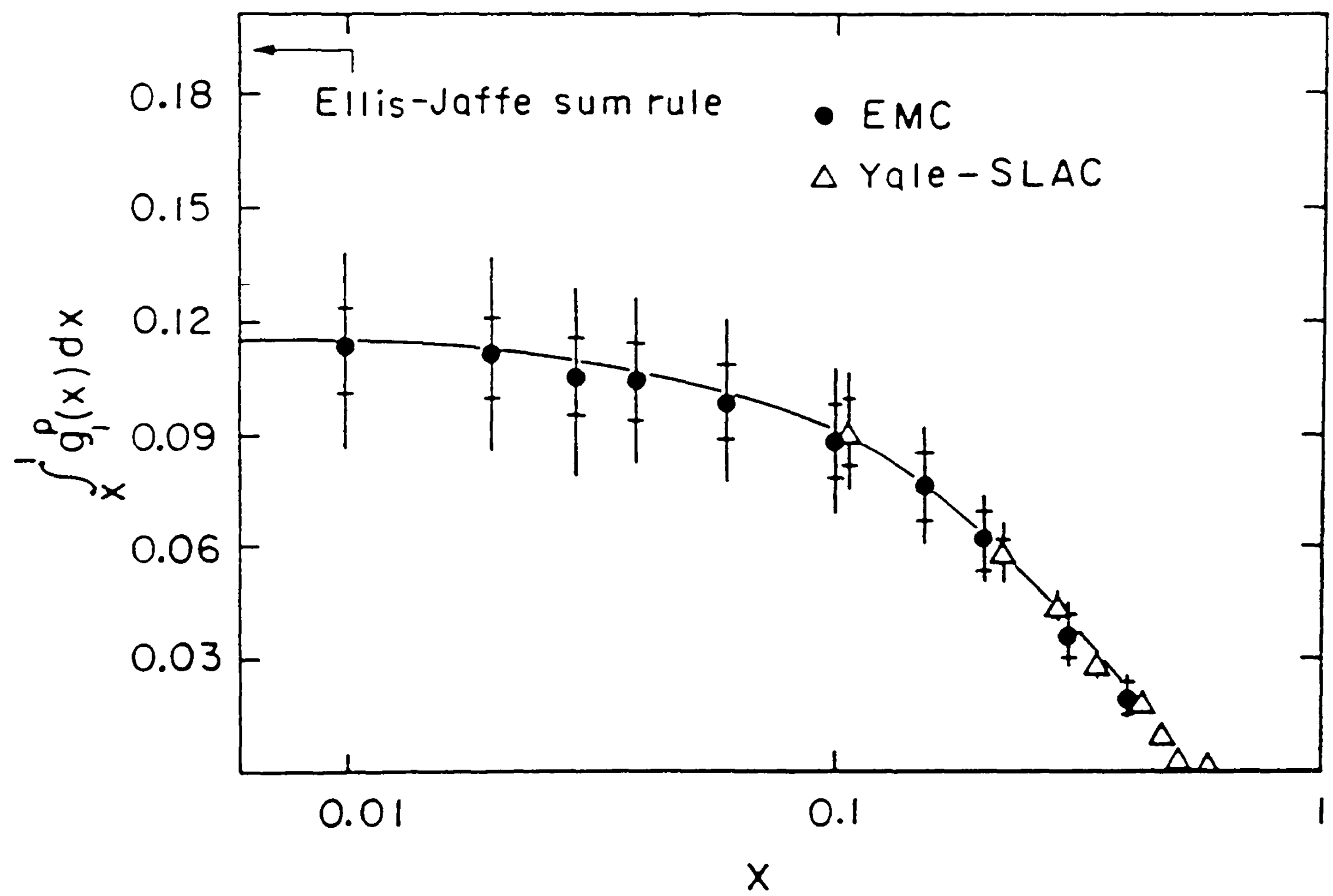
$$g_2(x, Q^2) = -g_1(x, Q^2) + \int_x^1 \frac{dy}{y} g_1(y, Q^2) + \bar{g}_2(x, Q^2) \tag{2.8}$$

where $\bar{g}_2(x, Q^2)$ should contain non-negligible quark mass effects and quark-gluon correlations. But it is this term which obscures the parton picture interpretation of g_2 [31,32].

Now we confine our attention to the interpretation of g_1. It is sufficient to know that, in the scattering of longitudinally polarised nucleons, the effects of g_2 in relation to g_1 are suppressed by a power of $\frac{\nu}{M^2}$.
The optical theorem states the proportionality of the total cross-section to the imaginary, i.e. absorptive, part of the scattering amplitude in forward direction. It is therefore clear, for the conservation of angular momentum, that a polarized photon, emitted by the left-handed μ^- detects the helicities of the partons inside a polarized proton.

The data from two experiments, from SLAC and from the EMC can be combined and the result is [4]:

$$\int_0^1 g_1^p(x, Q^2)dx = 0.126 \pm 0.018 \tag{2.9}$$



The average value of Q^2 for the EMC (SLAC) data is $\langle Q^2 \rangle = 10.7\ GeV^2$ $(\langle Q^2 \rangle = 5\ GeV^2)$. The actual value of Q^2 is different at each value of x where the asymmetry was counted. The x values were between 0.01 and 0.7, Q^2 varied from 1.5 to 70 GeV^2; $1.2 \cdot 10^6$ events were counted.

A necessary condition for the combination of SLAC and EMC data is the observed scaling behaviour of A which also suggests the interpretation of the asymmetry in the framework of the parton model.

It is important to note that the low transverse momentum of the partons with respect to the nucleon, makes them suitable for comparing current and constituent quarks. The direct physical interpretation of $\int_0^1 dx g_1(x,Q^2)$ can be given in terms of the quark and antiquark $\triangle q_i$ of flavour i and gluon $\triangle G$ contribution to the helicity of the proton:

$$\begin{aligned}\triangle q_i(Q^2) &\equiv \int_0^1 dx[q_i^{\uparrow}(x,Q^2)+\bar{q}_i^{\uparrow}(x,Q^2)-q_i^{\downarrow}(x,Q^2)-\bar{q}_i^{\downarrow}(x,Q^2)] \\ \triangle G(Q^2) &\equiv \int_0^1 dx[g^{\uparrow}(x,Q^2)-g^{\downarrow}(x,Q^2)]\end{aligned} \tag{2.10}$$

where $\uparrow$ and $\downarrow$ indicate parallel and antiparallel polarisation of the parton with respect to the proton polarisation, (i.e. $q^{\uparrow}(x),(g^{\uparrow}(x))$ is the probability of having a quark, (gluon), with momentum x and helicity $+1,(+\frac{1}{2})$ inside a proton of helicity $+\frac{1}{2}$).

At first sight one might think of the gluon contribution being suppressed by a small value of $\alpha_s(Q^2) \sim \frac{1}{lnQ^2}$. It was a first breakthrough in the spin problem when it was realised that the growth of $\triangle G(Q^2)$ with Q^2 compensates exactly the decrease of $\alpha_s(Q^2)$, so that the quantity $\alpha_s \triangle G$ is a constant to leading order [14,13].

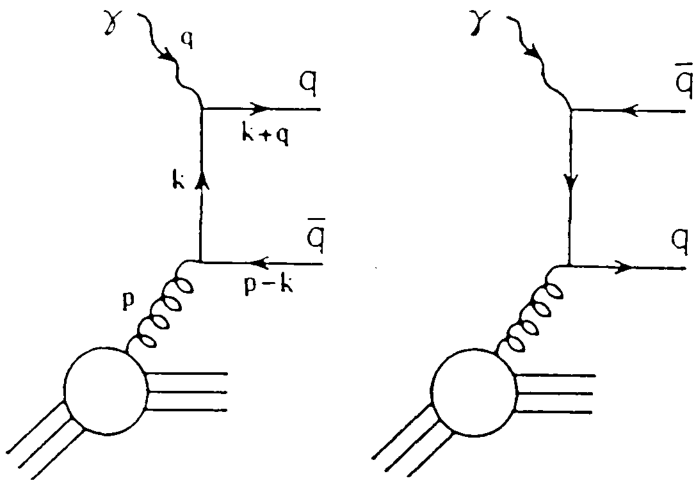
The next thing one has to know is the polarised photon parton cross-section. The quark photon cross-section is, of course, proportional to e_q^2, the quark contribution to $\int g_1$ being

$$\int_0^1 dx g_1(x,Q^2) = \frac{1}{2}\sum_i e_i^2 \triangle q^i(Q^2) \tag{2.11}$$

The calculation of the polarised photon gluon scattering process

$$\triangle\sigma^{\gamma g} = \sigma(\gamma^{\uparrow}g^{\uparrow}) - \sigma(\gamma^{\downarrow}g^{\uparrow}) \tag{2.12}$$

involves the evaluation of the following diagrams:



which has been done by several groups [13],[33,34],[35], with the result

$$\frac{d}{dcos\theta}[\triangle\sigma^{\gamma g}(x)] = \frac{\alpha}{2\pi}\sum_i \frac{e_i^2}{2}\left[x^2-(1-x)^2\right] \times\left[\frac{1}{1-\beta cos\theta}+\frac{1}{1+\beta cos\theta}-1\right] \tag{2.14}$$

where the sum is taken over all light quark flavours $i = u, d, s$. The infra-red singularity in this process has to be regularised either by taking the final state quarks off mass shell (i.e. mass regularisation)

$$\beta = 1 - \frac{2m^2 x}{Q^2(1-x)} \tag{2.15}$$

or by introducing a minimum angle or k_T transverse momentum of the produced outgoing q and $\bar{q}$ jets,

$$k_T^2 = \frac{Q^2}{4x}(1-x)sin^2\theta \tag{2.16}$$

The latter is also experimentally necessary in order to separate the two jets. The observability of the two jets is also the best criterion for the definition of $\triangle G$ in order to keep it clearly distinct from a quark distribution [36]. Also, having absorbed all infra-red sensitivity in the parton distribution, this definition of $\triangle G$ corresponds most closely to the intuitive constituent quark model picture of the nucleon: The only piece of the above process in which the virtual quark (or antiquark) is at low transverse momentum to the proton is the quark component.

Doing the θ integration, the partonic polarised gluon photon cross-section is:

$$\begin{aligned}\triangle\sigma^{\gamma g}(x) = \frac{\alpha_s}{2\pi}\sum_i \frac{e_i^2}{2}\left[x^2-(1-x)^2\right] \\ \times\left[ln\left(\frac{l^2}{Q^2}\right)+ln\left(\frac{x}{1-x}\right)\right]\end{aligned} \tag{2.17}$$

Neglecting the usual next to leading order QCD correction and multiplying with $\triangle G$, the gluon contribution to g_1 is:

$$\int_0^1 dx g_1(x) = -\frac{\alpha_s}{2\pi}\sum_i \frac{e_i^2}{2}\triangle G \equiv -\sum_i \frac{e_i^2}{2}\triangle\Gamma \tag{2.18}$$

Heavy quark effects are expected to have only minor (at most of order $\frac{\Lambda^2_{QCD}}{m_c^2}$) impact on $\int g_1$ as soon as the threshold for charm-production is passed. It has been suggested that in this case the intrinsic charm polarisation $\triangle c$ could even cancel the gluon part $\triangle\Gamma(m_c)$ [37]. It will be assumed to be negligible. (see also ref.[38])

There has also been some dispute over the sign of the gluon contribution to $\int g_1$ [13,15] vs.[14,5]. The fact that $\triangle G$ must be subtracted from $\triangle q$ is however easily seen by looking at the process in the Breitframe, neglecting intrinsic transverse momentum of the massless partons. Angular momentum conservation tells us that a left-handed muon emitting a left-handed photon can only be absorbed by a left-handed quark or antiquark. This left-handed quark or antiquark can emit a left-handed gluon which is however, after time reversal, a right-handed one. Therefore the detected helicities of (anti)quark and gluon are opposite [35].

Now we are ready to write down the partonic interpretation of $\int g_1$:

$$\int_0^1 dx g_1^p(x,Q^2) = \frac{1}{2}[\frac{4}{9}\triangle u(Q^2) + \frac{1}{9}\triangle d(Q^2) + \frac{1}{9}\triangle s(Q^2)] - \frac{1}{3}\triangle\Gamma(Q^2) \tag{2.19}$$

This can be further decomposed into:

$$\begin{aligned}\int_0^1 dx g_1^p(x,Q^2) = & \frac{1}{12}([\triangle u(Q^2) - \triangle d(Q^2)] \\ & + \frac{1}{3}[\triangle u(Q^2) + \triangle d(Q^2) - 2\triangle s(Q^2)] \\ & + \frac{4}{3}[\triangle u(Q^2) + \triangle d(Q^2) + \triangle s(Q^2) - 3\triangle\Gamma(Q^2)])\end{aligned} \tag{2.20}$$

This is the point where contact can be made with current algebra. The matrix element of the axial vector current

$$a_j s_\mu = < p|\bar{q}\gamma_\mu\gamma_5 \frac{\lambda_j}{2} q|p > \tag{2.21}$$

(where s_μ is the polarisation four vector), is protected from substantial QCD radiative corrections by the approximate conservation of A^j_μ in the flavour octet case. Because for $j \neq 0$, $dim[\partial_\mu A^{\mu j}] < 4$, the corresponding matrix elements of this 'soft' operator are independent of Q^2; A^j_μ has no anomalous dimension in this case. This is the reason why the Bjorken sum rule [39] is still valid, despite it predating QCD:

$$\int_0^1 dx(g_1^p(x) - g_1^n(x)) = \frac{1}{6} a_3 \tag{2.22}$$

Due to the isospin symmetry one just has to make the change $u \leftrightarrow d$ when considering the neutron structure function. The axial-vector isovector coupling of the nucleon

$$\frac{g_A}{g_V}\Big|_{n\to p} = a_3 = \triangle u - \triangle d \tag{2.23}$$

is measured in the neutron β-decay. The value of $\triangle u + \triangle d - 2\triangle s$ is more difficult to obtain from the data on semileptonic hyperon decays, because several assumptions about $SU(3)$ flavour symmetry have to be made. The explicit discussion of the matrix element a_8 and its extraction from hyperon decays will take place in the next chapter. Here for the sake of clarity, we adopt the approximate values to the most often cited values [40], viz

$$\triangle u - \triangle d \simeq 1.25 \qquad \text{and} \qquad \triangle u + \triangle d - 2\triangle s \simeq 0.6 \tag{2.24}$$

Neglecting also for the moment the small errors and putting these values back into equations (2.9) and (2.20), we find the surprising result:

$$\begin{aligned} 0.05 &= \triangle u(Q^2) + \triangle d(Q^2) + \triangle s(Q^2) - 3\triangle\Gamma(Q^2) \\ &= \triangle\Sigma - n_f \triangle\Gamma \end{aligned} \tag{2.25}$$

This astonishingly small number caused what is today known as the spin problem or the spin crisis [5]. The naive estimate of a_0 was made just by relating the matrix elements of the flavour singlet and flavour octet axial vector currents. Under the assumption that there are no strange quarks inside the proton Gourdin, Ellis and Jaffe [41] expected

$$\triangle u + \triangle d \simeq \triangle u + \triangle d - 2\triangle s \simeq 0.6 \tag{2.26}$$

$$\Rightarrow \qquad \int_0^1 dx g_1^p(x) = 0.187$$

One needed a huge gluon polarisation $\triangle G \sim 7$ and consequently a suspiciously well tuned opposite angular momentum in order to cure this problem exclusively with an additional gluon contribution [13]. This seems to be physically rather contrived because one would expect a dense spectrum of nearly degenerate baryon ground states, each with a different angular momentum.

One might think that other practical constraints on $\triangle G$ and $\triangle q$ are given by the unpolarised parton distributions

$$\begin{aligned} |q^\uparrow - q^\downarrow| &\leq q^\uparrow + q^\downarrow \\ |g^\uparrow - g^\downarrow| &\leq g^\uparrow + g^\downarrow \end{aligned} \tag{2.27}$$

However, these constraints have to be applied with special care because of the different scaling behaviour (i.e. different Q^2 dependence) of the unpolarised versus the polarised distributions. The moments of the polarised and unpolarised operators have different anomalous dimensions. If the inequalities are true at one scale Q_0^2, they are not necessarily true at another scale Q^2. Nevertheless, the comparison of the hypothetical polarised gluon distribution with existing data of unpolarised gluon contribution can be made under certain assumptions about the extrapolations into the extreme region $x \to 0$. Ross and Roberts [36] conclude that the best fit to the data can only be obtained after a negative strange quark polarisation $\triangle s < 0$ is added. Even if there has been evidence for $\triangle s = -0.15 \pm 0.08$ from elastic neutrino proton scattering from an earlier experiment [11], it needed to be confirmed by the

EMC measurement, because the latter required less model dependent analysis (see also chapter 5).

2.2. No quantitative prediction for a_0

A different problem arises with the question of how $\triangle\Sigma - n_f\triangle\Gamma$ relates to the flavour singlet axial vector current matrix element a_0. Because of the anomaly $dim[\partial^\mu A^0_\mu] = 4$, A^0_μ is not conserved even for massless quarks, and acquires an anomalous dimension γ_A. But the divergence of A^0_μ is proportional to the divergence of a different operator:

$$\partial^\mu A^o_\mu = n_f \frac{\alpha_s}{2\pi} trGG^* = n_f \partial^\mu K_\mu \tag{2.28}$$

where

$$K_\mu = \frac{\alpha_s}{2\pi} \epsilon_{\mu\nu\lambda\sigma} tr[A^\nu(G^{\lambda\sigma} - \frac{2}{3}A^\lambda A^\sigma)] \tag{2.29}$$

is clearly gauge-dependent. But $A^0_\mu - n_f K_\mu$ is conserved for massless quarks, and it has been argued [38] that one should identify $< p|A^0_\mu - n_f K_\mu|p >$ with $\triangle\Sigma$ in the parton language. This identification is legitimate even though the quantities are not necessarily the same, but because their anomalous dimensions are identical. The same is true for the identification of $\triangle\Gamma$ and K_μ. The gauge dependent operator K_μ actually has gauge invariant diagonal matrix elements, at least for gauge transformations which do not change the winding number $n \propto \int d^4x TrGG^*$. Mandula [42] calculated $< p|K_\mu|p >$ on the lattice using the above identification and found that the contribution of $-\frac{3\alpha_s}{2\pi}\triangle G$ is less than 0.05 in magnitude.

The suggested decomposition of the flavour singlet a_0, motivated by the parton picture, is:

$$a_0|_{Q^2} = \triangle q|_{Q^2} - n_f \triangle\Gamma|_{Q^2} \tag{2.30}$$

The non-vanishing anomalous dimension γ_A controls via the RG equation the Q^2 dependence:

$$a^o(Q^2) = a^o(\mu^2) exp[\int_{\bar{g}(Q^2)}^{\bar{g}(\mu^2)} d\bar{g} \frac{\gamma_A(\bar{g})}{\beta(\bar{g})}] \tag{2.31}$$

where $\bar{g}(Q^2)$ is the effective coupling, and

$$Q^2 \frac{d}{dQ^2} \bar{g}(Q^2) = \beta(\bar{g}) \tag{2.32}$$

is the QCD RG β-function. $\gamma_A(\bar{g}^2)$ which begins at two loop order has been calculated perturbatively [43]. The resulting value

$$\gamma_A(\bar{g}) = \left(\frac{\bar{g}^2}{16\pi^2}\right)^2 16 n_f + \ldots \tag{2.33}$$

is actually too small for explaining the discrepancy between naive expectation $a_0 \sim a_8$ and experiment $a_0 \sim 0$. The corrections to a_0 are of order $\alpha_s(Q^2)$ and the QCD corrected value is

$$\begin{aligned} a^o(Q^2) &\sim (1 + \frac{1}{\pi} \frac{6n_f}{33 - 2n_f} [\alpha_s(Q^2) - \alpha_s(\mu^2)]) a^o(\mu^2) \\ Q^2 &\to \infty \end{aligned} \tag{2.34}$$

The calculation of γ_A depends on the assumption of the validity of perturbative QCD in its extrapolation to low $Q^2 \to \mu^2$. These extrapolations work perfectly well in the flavour octet case. Nevertheless, for a critique of the calculation in the flavour singlet case the reader is referred to Ball's attempt to calculate γ_A non-perturbatively [44].

In general, a quantitative prediction of $a_0(0)$ 'still appears to be beyond the reach of current techniques in non-perturbative QCD'. This is the conclusion of a recent exhaustive discussion by Shore and Veneziano [17] who study the equivalent of the Goldberger-Treiman (GT) [45] relation for the U(1) channel. The GT relation

$$f_\pi g_{\pi NN} = M g_A \tag{2.35}$$

where $g_{\pi NN}$ is the pion nucleon coupling constant, has a firm theoretical foundation in terms of QCD and SSB. It is an exact identity in the chiral limit. Because of the RG dependence of a_0 and because the η' is not a Nambu-Goldstone (NG) boson, the appropriate form of the U(1) GT relation can only be written in terms of an unphysical 'would be' NG boson of the U(1) channel:

$$f_{\eta_o}(0) g_{\eta_o NN}(0) = 2M a_o(0) \tag{2.36}$$

with the RG invariant factor $f_{\eta_0}(0)$ whose square is proportional to the first momentum of the topological susceptibility in QCD:

$$f_{\eta_o}(0) = 2n_f \left(\frac{\partial}{\partial k^2} \int dx e^{ikx} i < 0|T^* QQ|0 > \right)^{\frac{1}{2}} \Big|_{k=0} \tag{2.37}$$

$$\text{with} \qquad Q = \frac{\alpha_s}{8\pi} trGG^*$$

So far there exists no estimate of this quantity, and one might expect it to be small in order to make a_0 small. On the other hand one might argue:

$$g_{\eta_o NN} = 0 \tag{2.38}$$

This is true in any model where there is no U(1) component; e.g. in the Skyrme picture the nucleon is the soliton constructed from Goldstone fields of the coset space $SU(3)_L \otimes SU(3)_R/SU(3)$ [46] and equation (2.38) follows easily [47].
The above GT like relation can also be rewritten in terms of η':

$$F g_{\eta' NN} + \frac{1}{n_f} F^2 m_{\eta'}^2 g_{GNN}(0) = 2M a_o(0) \tag{2.39}$$

and given the additional smoothness assumption

$$g_{\eta' NN}(m_{\eta'}^2) \simeq g_{\eta' NN}(0) \tag{2.40}$$

F is the fundamental constant, as opposed to the incorrect (RG variant) '$f_{\eta'}$' decay parameter, and g_{GNN} stands for a glue-nucleon coupling whose RG properties are of a highly non-trivial nature. Unfortunately, it is not possible to measure the universal constant F, because each process in which it appears involves a new and unknown proper vertex $g_{GXX}(0)$. For example, in the relation describing the process $\eta' \to \gamma\gamma$ a glue-photon vertex has to be taken into account which is equally unknown:

$$F g_{\eta'\gamma\gamma} + \frac{1}{n_f} F^2 m_{\eta'}^2 g_{G\gamma\gamma}(0) = \frac{4N_c}{3\pi}\alpha_{em} \tag{2.41}$$

again with the additional smoothness assumption

$$g_{\eta'\gamma\gamma} \equiv g_{\eta'\gamma\gamma}(m_{\eta'}^2) \simeq g_{\eta'\gamma\gamma}(0) \tag{2.42}$$

The term on the right-hand side of equation (2.41) is due to the QED axial anomaly and occurs in the simpler version of the $\pi^0 \to \gamma\gamma$ decay as well:

$$f_\pi g_{\pi\gamma\gamma} = \frac{N_c}{3\pi}\alpha_{em} \tag{2.43}$$

But again, due to the QCD axial anomaly, it is not possible to generalise equation (2.43) by replacing f_π with '$f_{\eta'}$'.

With this important analysis of the $\eta' \to \gamma\gamma$ decay and the GT U(1) relation Shore and Veneziano clarified the criticism of incorrectly relating a_0 to '$f_{\eta'}$. The unfortunate consequence of this result is of course that, in the foreseeable future, no reliable calculation of a_0 can be performed, and that one has to look for other ways to solve the spin problem.

3. Semileptonic Baryon Decays

3.1. SU(3) breaking in hyperon β-decays

Our understanding of the β-decays in the standard model is based on the observation that the quark mass eigenstates are not the same as the weak eigenstates. All elementary fermion fields ψ_i which play a role in weak currents are left-handed and transform as doublets under SU(2):

$$\psi_i = \begin{pmatrix} \nu_i \\ l_i^- \end{pmatrix} \qquad \text{and} \qquad \begin{pmatrix} u_i \\ d_i' \end{pmatrix} \tag{3.1}$$

where $i = 1,2,3$ indicates the fermion family. By convention, the three charge 2/3 quarks (u,c,t) are unmixed, and all the mixing between mass and weak eigenstates is expressed in terms of the 3×3 unitary Kobayashi-Maskawa matrix V [48,1] operating on the charge -1/3 quarks (d,s,b).

$$\begin{pmatrix} d' \\ s' \\ b' \end{pmatrix} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \begin{pmatrix} d \\ s \\ b \end{pmatrix} \tag{3.2}$$

The Kobayashi-Maskawa matrix is a generalisation of the four quark case, where the matrix can be parametrised by the single Cabibbo angle θ_C [49]. It is essential to test the unitarity of this matrix

$$|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 1 \qquad \text{etc.} \tag{3.3}$$

In order to evaluate the matrix elements V_{ud} and V_{us} from the β-decays of the flavour octet baryons one has to understand precisely if and how SU(3) symmetry breaking effects alter the necessary SU(3) parametrisation of these decays. The matrix element $\mathcal{M}$ for the decays is given by

$$\mathcal{M} = \frac{G_F}{\sqrt{2}} < B'|J^\mu|B > \bar{u}_l(p_l)\gamma_\mu(1-\gamma_5)u_{\bar{\nu}}(p_{\bar{\nu}}) \tag{3.4}$$

where $G_F = 1.116 \cdot 10^{-5} GeV^{-2}$ is the universal weak coupling constant ($\hbar c = 1$), and p_l and $p_{\bar{\nu}}$ are the lepton and anti-neutrino four-momenta. The baryon term can be decomposed into its most general form as follows:

$$\begin{aligned} < B'|J^\mu|B > = C\bar{u}_{B'}(p')\{f_1(q^2)\gamma^\mu + i\frac{f_2(q^2)}{M}\sigma^{\mu\nu}q_\nu + \frac{f_3(q^2)}{M}q^\mu \\ + [g_1(q^2)\gamma^\mu + i\frac{g_2(q^2)}{M}\sigma^{\mu\nu}q_\nu + \frac{g_3(q^2)}{M}q^\mu]\gamma_5\}u_B(p) \end{aligned} \tag{3.5}$$

where $C = [V_{ud}, V_{us}] \simeq [cos\theta_c, sin\theta_c]$ for $\triangle S = [0,1]$ strangeness conserving or changing transitions, θ_c is the Cabibbo angle, and $q = p - p'$, where p (p') and M (M') are momenta and masses of the B (B') baryons respectively.

Time reversal invariance implies that f_i and g_i are real. Contributions to the decay distributions from f_3 and g_3 are proportional to the e^- mass divided by the baryon mass and can therefore be neglected. The contributions of f_2 and g_2 to the decay amplitudes is suppressed by the momentum transfer and therefore vanish in the SU(3) symmetry limit of degenerate baryon masses. With the hypotheses of the Cabibbo theory [49], the scalar matrix elements of the hadronic weak current at zero momentum transfer are completely determined by two parameters. These define the strength of the antisymmetric F and symmetric D couplings of two octets to form a third octet, and by the Cabibbo angle θ_c.

The construction of traceless baryon and axial vector matrices $\mathcal{B}$ and $\mathcal{A}$ is therefore necessary, so that the matrix element may be expressed in the frequently used form:

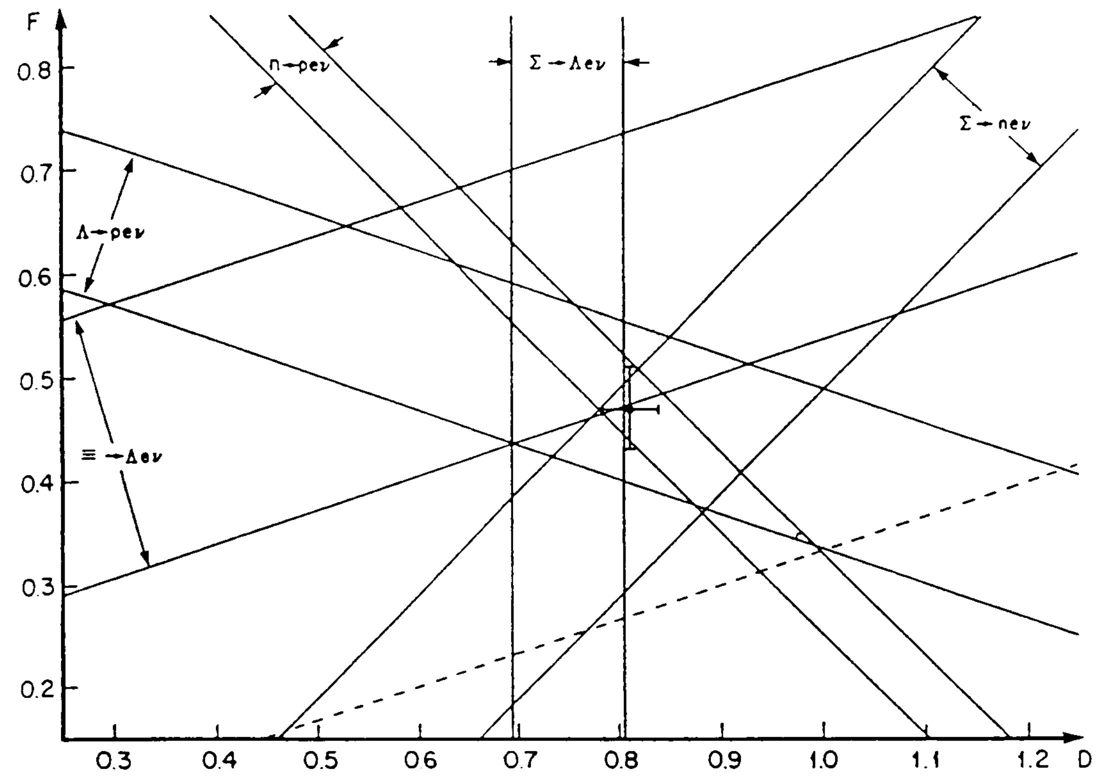
$$< B'|A^{\mu}|B >= Dtr(\mathcal{A}\{\mathcal{B},\bar{\mathcal{B}}\}) + Ftr(\mathcal{A}[\mathcal{B},\bar{\mathcal{B}}]) \tag{3.6}$$

These two parameters can be fixed on the basis of two values for $\frac{g_1}{f_1}|_{B\rightarrow B'}$ taken for two different baryon decays. The first and most precisely determined input is the value of the neutron decay, which is taken from the Particle Data Group [50]. It is a little larger than the 'old' and often cited (e.g. [40]) value 1.259 ± 0.004:

$$\frac{g_1(0)}{f_1(0)}|_{n\rightarrow p} := \frac{g_A}{g_V}|_{n\rightarrow p} = 1.261 \pm 0.004 = F + D \tag{3.7}$$

The real problem is the second input because of possible SU(3) symmetry breaking effects. The Ademollo-Gatto theorem [51] tells us that SU(3) breaking enters g_V at second order only and hence with minor consequences. But g_A is not similarly protected against any corrections stemming from the non-degenerate baryon mass differences $M_B - M_{B'}$.

Normally, the resulting F/D value is extracted from a least squares fit to all β-decays. In the graph below the result of the analysis of Jaffe and Manohar [40] is shown, indicating the 'allowed' regions of F and D for various hyperon decays under the assumption $g_2 = 0$, and marking the point which is the best fit value:



Their value for F/D

$$F/D = 0.58 \pm 0.05 \tag{3.8}$$

is calculated without including any SU(3) breaking effects at all. Naturally therefore Jaffe and Manohar assume a large error 'even if all hyperon β-decay measurements were consistent within very small errors with a single choice of F and D'.

The dashed line in the figure above marks the points where one can have no strange quarks in the proton and agreement with the EMC result simultaneously without a gluon contribution. This line is two standard deviations away from the 'best fit', i.e. from the already very low value of F/D. In the light of the proton spin structure measurement the analysis of F/D is very important and will be discussed explicitly below, after a new value of F/D has been found.

Until today no general consensus exists about the pattern in which SU(3) is broken and how it would change the F/D analysis. The relatively poor fit in the SU(3) limit shows, however, that SU(3) symmetry breaking in hyperon decays is present. In the following we shall study a few proposed patterns of SU(3) symmetry breaking.

For example one of the most recent analysis is Zenczykowski's discussion [52] of two kinds of potential quark-model effects which invalidate SU(3) symmetry parametrisation due to wave-function deformations and configuration mixing. In the framework of a harmonic oscillator quark model of baryons with the groundstate wavefunction ψ_0, the presence of quarks with unequal masses gives rise to two different frequencies α_ρ and α_λ:

$$\psi_0 = \left(\frac{\alpha_\rho \alpha_\lambda}{\pi}\right)^{\frac{3}{2}} e^{-\frac{1}{2}\left(\alpha_\rho^2 \rho^2 + \alpha_\lambda^2 \lambda^2\right)} \tag{3.9}$$

where $\rho = (r_1 - r_2)/\sqrt{2}$ and $\lambda = (r_1 + r_2 - 2r_3)/\sqrt{6}$ are the two relative coordinates. As a result of a possible mismatch between the wavefunctions of the initial and final baryons the form factors are subject to a modification by a factor of $O(exp[-q^2/6\alpha^2])$, where

$$q^2 = (M' - M)^2 \left[1 + \left(\frac{M' - M}{2M}\right)\right]^2 \tag{3.10}$$

is the generally nonzero three-momentum transfer. In addition to this correction the necessary relativistic corrections due to the small components of the Dirac spinors are quark mass dependent and break therefore SU(3) symmetry also.

Taking into account all these corrections, Zenczykowski [52] calculated the deviations of the formfactors from their NQM value of g_A/g_V, under the assumption that the SU(3) breaking parameter $\delta \equiv (m_s - m_u)/m_u$, (i.e. the ratio of the constituent quark masses), is not too large:

$$
\begin{array}{ll}
n \to p: & \frac{g_A}{g_V} = \frac{g_A}{g_V}\Big|_{NQM}[1 - \frac{\alpha^2}{3m_u^2}] \\
\Lambda \to p,\ \Sigma^- \to n: & \frac{g_A}{g_V} = \frac{g_A}{g_V}\Big|_{NQM}[1 - (1 - \frac{5\delta}{6})\frac{\alpha^2}{3m_u^2}] \\
\Xi^- \to \Lambda, \Sigma^o or\ \Xi^o \to \Sigma^+: & \frac{g_A}{g_V} = \frac{g_A}{g_V}\Big|_{NQM}[1 - (1 - \frac{3\delta}{4})\frac{\alpha^2}{3m_u^2}] \\
\Xi^- \to \Xi^o: & \frac{g_A}{g_V} = \frac{g_A}{g_V}\Big|_{NQM}[1 - (1 + \frac{\delta}{6})\frac{\alpha^2}{3m_u^2}]
\end{array}
\qquad (3.11)
$$

He concluded that the SU(3) breaking induced by the mass terms works in the wrong direction. This is because the relativistic corrections are of course smaller for the heavier strange quarks $\delta \geq 0$ than for the light up and down quarks, and is in contradiction with the expectation that a large δ induces large SU(3) breaking.

Another way to modify the SU(3) invariant description of semileptonic baryon decays with the F and D parameters lies in the possible generalisation of the groundstate baryon wavefunctions. Isgur and Karl [53] originally explained this configuration mixing (between N, Λ, Σ, Ξ states of the $(8_S; 56, 0^+), (8'_S; 56', 0^+)$ and $(1_M, 8_M, 10_M; 70, 0^+)$ multiplets of the symmetric quark model) as a consequence of SU(6) breaking colour-hyperfine interactions in the framework of the harmonic oscillator quark model. Due to this freedom of choice in different parameters (e.g. the strength of hyperfine interactions) the calculated mixing angles are not unique and Zenczykowski [52] cites them as 'tailored to our needs'. By including mass independent relativistic corrections with an SU(3) invariant scaling factor, and by neglecting D-wave 'impurities' Zenczykowski also calculated the modified form factors starting with the g_A/g_V ratios of the NQM. The resulting F and D parameters for this case, turn out to be meaningful only to an 15% accuracy level. (Table III in ref.[52])

TABLE III. Ratios of axial-vector-to-vector couplings g_A/g_V modified by configuration mixing. The underlined entry constitutes a fit to the experimental number.

Process	$(g_A/g_V)^{SU(6)}$	g_A/g_V modified	+ relativistic effects	Experiment	SU(3) param.
$n \to pe^-\nu$	5/3	1.60	<u>1.258</u>	1.258 ∓ 0.004	$F+D$
$\Sigma^- \to ne^-\nu$	$-1/3$	-0.32	-0.25	$\mp(0.36 \mp 0.04)$	$F-D$
$\Lambda \to pe^-\nu$	1	0.93	0.73	0.696 ∓ 0.025	$F+\frac{D}{3}$
$\Xi^- \to \Lambda e^-\nu$	1/3	0.35	0.28	0.25 ∓ 0.05	$F-\frac{D}{3}$
$\Xi^- \to \Sigma^0 e^-\nu$	5/3	1.70	1.34		$F+D$

The possibility of a wavefunction overlap due to a different spatial dependence for u, d, and s quarks has also been discussed by Donoghue, Holstein and Klimt (DHK) [54] in the framework of a bag model approach to hadron structure. In addition to this correction these authors discussed the kinematic 'center of mass' or 'recoil' corrections. By superposing several momentum eigenstates of the baryons with some normalised weighting factor $\phi_B(p)$

$$|\psi_B>= \int d^3p\phi_B(p)|\psi_B(p)> \tag{3.13}$$

an estimate (in terms of the free parameter $< p^2 >$) was given for the change in form factors due to the recoil of the baryons.

$$\begin{aligned}\int d^3x < \Psi_{B'}|A^3_\mu|\psi_B > &= g_A^{B'B} \int d^3p\phi^*_{B'}(p)\phi_B(p) \\ &\quad (2\pi)^3\bar{\psi}_{B'}(p)\gamma_\mu\gamma_5\frac{\lambda^3}{2}\psi_B(p) \\ &\equiv g_A^{B'B}\rho_A^{B'B}\end{aligned} \tag{3.14}$$

$$\text{with} \quad \rho_A^{B'B} = 1 - \frac{< p^2 >_{B'B}}{3M'M}[\frac{1}{4} + \frac{3}{8}\frac{M'}{M} + \frac{3}{8}\frac{M}{M'}]$$

The square momentum $< p^2 > \neq 0$ can be taken from a bag model calculation or can be fit to all data [55]. Because in the case of g_1 both SU(3) breaking effects (wavefunction overlap and recoil) go in opposite directions, the overall amount of SU(3) breaking is less than one might have expected. Thus this ansatz alone cannot explain the observed magnitude of SU(3) symmetry breaking.

Another result of the DHK investigation has been the calculated shift in V_{us} which after the inclusion of SU(3) breaking amounts only to 2% . The more exact value of V_{us} is still determined by K_{e3} decays [50,56]

$$K^+ \to \pi^o e^+ \nu_e, \; K^+ \to \pi^o \mu^+ \nu_\mu, \; K_L^o \to \pi^- e^+ \nu_e, \; K_L^o \to \pi^- \mu^+ \nu_\mu \qquad (3.15)$$

These decays are not plagued by the theoretical uncertainty of the axial vector current matrix elements, because only the vector current contributes:

$$\mathcal{M} \propto V_{us} < \pi, p | \bar{u} \gamma_\mu s | K, p > \qquad (3.16)$$

Of course, common to all evaluations of the matrix elements V_{ij} is the calculation of the radiative corrections to the process in question, and subject to further but small theoretical uncertainty.

More important for our discussion than the accurate convergence of the V_{ij} matrix, is the fact that all the above phenomena, which are mostly modifications of the NQM, should match the description of the baryon magnetic moments. Until now however, it is not known why the magnetic moments seem to violate SU(3) much more strongly. It is one goal of this thesis to point out that there is still room for the existence of a simple but consistent pattern in which the quark mass difference gives rise to a significant shift of the old hyperon values of g_A/g_V, leading to a new

F/D value. This new value will also fulfill the requirement of being derivable from the baryon magnetic moments.

This still leaves open the question why and how different quark masses can shift the $\frac{g_A}{g_V}$ ratios, after they have been measured with incredible precision? This will be addressed in the next section.

3.2. The 'induced second class' form factor and a new F/D

A glance at the entries of $\frac{g_A}{g_V}$ in the Particle Data Book [50] shows that these ratios have been almost exclusively analysed under the *theoretical* assumption $g_2 = 0$, which a *priori* is only valid in the perfect SU(3) limit.

The most recent data [57] for the axial-vector-to-vector form factor ratio $\frac{g_1(0)}{f_1(0)}$ of the $\Sigma^- \to n e^- \bar{\nu}$ decay, however, included for the first time an independent extraction of the induced pseudo-tensor-to-vector form factor ratio $\frac{g_2(0)}{f_1(0)}$. In previous experiments [58] the statistics were not sufficient in order to measure the ratio $\frac{g_1}{f_1}$ without the assumption $g_2 = 0$. Without this constraint the new experimental result is a reduced value for $\frac{g_1(0)}{f_1(0)}$. Instead of $\frac{g_1(0)}{f_1(0)} = -0.328 \pm 0.019$ and $g_2(0) = 0$ fixed, with the least square value $\chi^2/N_{DF} = 2.52/3$ for the fit, one has now with $\chi^2/N_{DF} = 1.07/2$:

$$\frac{g_1(0)}{f_1(0)}\Big|_{\Sigma^- \to n} = -0.20 \pm 0.08 \qquad \text{and} \qquad \frac{g_2(0)}{f_1(0)}\Big|_{\Sigma^- \to n} = 0.56 \pm 0.37 \tag{3.17}$$

The ratio $\frac{g_1}{f_1}$ of all other hyperon decays could only be determined so far under the assumption $g_2 = 0$.

The invariance of strong interactions under the G-parity operation $G = Ce^{i\pi I_2}$, the product of charge symmetry and charge conjugation allows us to classify the terms which appear in the baryonic weak matrix element according to their transformation property under G:

$$\begin{aligned} GJ_{i1}G^{-1} &= +J_{i1}\cdot\xi \\ GJ_{i2}G^{-1} &= -J_{i2}\cdot\xi \end{aligned} \tag{3.18}$$

where $\xi = \pm 1$, depending on f_i or g_i.

After it has been established that the standard left-handed quark current, classified as first-class current, is the only weak interaction current, the so-called second-class currents (which have opposite G-parity to the first-class currents) have to vanish [59]. This is the reason why the term proportional to g_2 is normally set to zero. g_2 is called the 'induced second-class' or the 'weak electricity' form factor. Its observed appearance in Σ^--decay, however, does not cast any doubt on the fact that the weak hadronic current is a left-handed quark current. The g_2 term is expected to be non-zero, given that SU(3) is broken by the difference between the mass of the strange-quark and the mass of the up- and down-quark. This does not however, imply the existence of second class currents.

The invariance under the isospin transformation in (3.18) is obviously not altered by a different strange quark mass. But the generalisation of the G-parity operation to $G' = Ce^{i\pi V_2}$-parity operation involves the V-spin symmetry, which is broken by the unequal quark masses. Therefore, strong interaction effects can be expected to generate a non-zero g_2 term in a strangeness changing transition of the order

$$\frac{E_l}{M+M'} \quad , E_l: \ \text{lepton energy} \tag{3.19}$$

Only in the β-decay of the neutron is this term definitely negligible.

In the NQM and the MIT bag models g_2 can be calculated, because it is related to the electric dipole moment

$$\vec{d}_{ij} = -i \int d^3x \vec{r} A^o_{ij}(x) = \frac{1}{2}\left[\frac{1}{2m_j} - \frac{1}{2m_i}\right] \vec{\sigma}_{ij} \tag{3.20}$$

where A^0 is the time component of the axial vector current $A^{\mu}_{ij} = \psi_i \gamma^{\mu} \gamma_5 \psi_j$, m_i, m_j are the quark masses and $\vec{\sigma}_{ij}$ is the matrix element of the spin operator connecting quarks i and j between appropriate baryon states $|B>$.

After introducing the relativistic corrections as an overall normalisation due to the reduced $\frac{g_A}{g_V}|_{n\to p}$ ratio from 5/3 to 5/4, Donoghue and Holstein [60] calculated for the strangeness changing process g_2 to be positive, $g_2 > 0$, and they found the astonishingly high value of $\frac{g_2}{g_1}|_{NQM} \approx 0.73$ in the case of the Λ-particle β-decay and values of similar magnitude for the other $\triangle S = 1$ processes.

The same authors calculated $\frac{g_2}{g_1}$ in the framework of the relativistic MIT bag model, where its value $\frac{g_2}{g_1}|_{MIT} \approx 0.30$ is half of its NQM value.

That confinement effects are responsible for the suppression of the induced second-class form factors has been argued by Carson, Oakes and Willcox [61] who perform an extensive analysis in the framework of a modified bag model by taking into account all kinds of corrections. These concern the kinematics and even the adaptation of a perturbative QCD motivated renormalisation. Their result is that g_2 turns out to be negligible.

Pointing out the spurious reference-frame dependence of four different bag model calculations, Lie-Svendson and Høgaasen [62] attempt to restore the Lorentz invariance by including recoil effects to the boosted three quark baryon system. Even if Lorentz-invariance cannot be completely satisfied, these authors claim that frame dependence can be reduced to an acceptable level. For this purpose all calculations have been performed in two different frames. The first frame is the center of mass frame in which the initial baryon is at rest, the second is a modified version of the

Breit-frame in which the two baryons, instead of having opposite momenta, have opposite velocities. Most interesting is the drastic reduction of g_2 from its value in the static limit as soon as the recoil effect is respected. Another interesting reduction of g_1 and g_2 happens in versions of the bag model where pseudoscalar meson fields are taken into account. The corresponding calculations introduce unwanted pion pole contributions to g_1 and g_2 which have consequently to be regularised in an appropriate manner. The conclusion of Lie-Svendson and Høgaasen concerning g_2 is that no bag model can reproduce a value of g_2 as large as the one observed (3.17). However, the sign of g_2 they find is at least in agreement with the data. Another interesting fact is their calculation of the $\frac{g_1}{f_1}$ ratios. Especially for the chiral bag model [63],[64],[65] which respects chiral symmetry unlike other bag models the following values have been calculated (from Table IV in [62]):

$$
\begin{aligned}
n \to p : \quad & \left.\frac{g_1}{f_1}\right|_{n\to p} = 1.26 \ , \quad \Sigma^- \to n : \quad \left.\frac{g_1}{f_1}\right|_{\Sigma^-\to n} = -0.24 \\
\Lambda \to p : \quad & \left.\frac{g_1}{f_1}\right|_{\Lambda\to p} = 0.76 \ , \quad \Xi^- \to \Lambda : \quad \left.\frac{g_1}{f_1}\right|_{\Xi^-\to \Lambda} = 0.25 \qquad (3.21) \\
\Sigma^- \to \Lambda : \quad & \left.\frac{g_1}{f_1}\right|_{\Sigma^-\to \Lambda} = 0.59 \ , \quad \Xi^- \to \Sigma^o : \quad \left.\frac{g_1}{f_1}\right|_{\Xi^-\to \Sigma^o} = 1.25
\end{aligned}
$$

The question remains open, as to whether future experiments can confirm the data on g_2 by measuring other β-hyperon decays with similar precision and a subsequent 'best fit analysis'. But because these data from the Fermilab experiment already exist, it is legitimate to argue again for an SU(3) symmetry of the axial vector currents.

The mass differences at the current quark level do not necessarily invalidate the SU(3) symmetry in the polarisation of the baryonic wavefunctions, or Cabibbo's assumption of a valid F and D parametrisation of the hadronic axial vector current.

However, the SU(3) symmetry breaking due to unequal quark masses does take place and is observed through a non-vanishing g_2 form factor in semileptonic hyperon decays.

In order to determine the corresponding F/D value, the following ansatz is chosen, which is SU(3) symmetric in the polarization but not in the quark-masses:

$$\triangle s \cdot s_\mu =< p|\bar{s}\gamma_5\gamma_\mu s|p >=< \Sigma^+|\bar{d}\gamma_5\gamma_\mu d|\Sigma^+ >=< \Xi^-|\bar{u}\gamma_5\gamma_\mu u|\Xi^- > \tag{3.22}$$

with p, Σ^+, Ξ^- being in the same spin eigenstate. Given that the F/D ratio is an invariant characteristic of the SU(3) symmetric axial-vector matrix-elements, it must be taken seriously. Nevertheless, for practical reasons and for the sake of clarity, the following ratio between the matrix elements of the two axial-vector diagonal flavour octet currents is taken for the analysis

$$z = \frac{a_8}{\sqrt{3}a_3} = \frac{F/D - \frac{1}{3}}{F/D + 1} \tag{3.23}$$

The measured values are [50,57]

$$\left.\frac{g_1(0)}{f_1(0)}\right|_{n\to p} = \triangle d(n) - \triangle d(p) = \triangle u - \triangle d = a_3 = 1.261 \pm 0.004 \tag{3.24}$$

and

$$\left.\frac{g_1(0)}{f_1(0)}\right|_{\Sigma^-\to n} = \triangle s(\Sigma^-) - \triangle s(n) = \frac{1}{2}a_3(3z-1) = -0.20 \pm 0.08 \tag{3.25}$$

$$[-0.328 \pm 0.019 \ , \ g_2 = 0]$$

The value with the constraint $g_2 = 0$ is shown in brackets, and $\triangle q$ is twice the helicity expectation value of the current quark q, (i.e. the fraction of the baryon spin carried by q), evaluated inside a polarized proton or the indicated baryon. The result

$$z = 0.23 \pm 0.04 \quad [0.160 \pm 0.010] \tag{3.26}$$

or equivalently,

$$F/D = 0.73 \pm 0.09 \quad [0.59 \pm 0.02] \tag{3.27}$$

shows an increase of z and F/D as soon as the constraint $g_2 = 0$ is relaxed. The value in brackets is obtained if $g_2 = 0$. Of course, the error is still very large. $z = 0.2$ corresponds to the SU(6) value $F/D = 2/3$. Not only in restriction to the NQM this value may indeed be the best criterion for a flavour-spin-symmetric groundstate.

The ratios $\frac{g_1(0)}{f_1(0)}$ for the other hyperon β-decays, taking into account possible effects of the difference in the constituent quark masses, i.e. taking into account g_2, can of course be easily predicted by fixing $a_3 = 1.261 \pm 0.004$, and by taking $z = 0.23 \pm 0.04$ $[z = 0.160 \pm 0.010]$:

$$\frac{g_1(0)}{f_1(0)}\Big|_{\Xi^- \to \Lambda} = \triangle s(\Xi^-) - \triangle s(\Lambda) = z a_3 = 0.29 \pm 0.05 \tag{3.28}$$

$$[0.202 \pm 0.013]$$

$$\frac{g_1(0)}{f_1(0)}\Big|_{\Lambda \to p} = \triangle s(\Lambda) - \triangle s(p) = \frac{1}{2} a_3 (1 + z) = 0.77 \pm 0.03 \tag{3.29}$$

$$[0.731 \pm 0.006]$$

These different predictions should be compared with the experimental data [50] for both the Ξ^- and the Λ decay which unfortunately are only available under the constraint $g_2 = 0$:

$$\left.\frac{g_1(0)}{f_1(0)}\right|_{\Xi^- \to \Lambda} = 0.25 \pm 0.05 \ , \ \text{with } g_2 = 0 \tag{3.28'}$$

$$\left.\frac{g_1(0)}{f_1(0)}\right|_{\Lambda \to p} = 0.718 \pm 0.015 \ , \ \text{with } g_2 = 0 \tag{3.29'}$$

There is good reason to expect similar corrections for the form factor ratios $\frac{g_1(0)}{f_1(0)}$ (due to a non-vanishing g_2) for all the semileptonic hyperon decays, in which the V-spin is the underlying symmetry. The corrections for $\frac{g_1}{f_1}$ (with respect to a non-vanishing g_2) must have the same sign. The dominant modification of the form factors is described by the linear combination $g_1 - \epsilon g_2$ with $\epsilon = (m_Y - m_B)/m_Y$ for the hyperon β-decay $Y \to B$, with a positive value for g_2. At the moment the expected changes ($3.28' \to 3.28$) and ($3.29' \to 3.29$) are consistent with the data (3.25). Because of the large errors in the measurement of g_2 (3.17) and z (3.26) this is not such a surprise. Numerically, the old predictions in square-brackets (3.28/29) are not yet excluded.

It is possible that in the near future, high precision measurements at FERMILAB [66] allow to determine $\frac{g_2}{f_1}$ in the β-decay $\Lambda \to p$. Thus, the new value of z (or F/D) may be confirmed, with a consequently much smaller error. At present it looks as if such a new value will be significantly distinct from the old one.

If future data however support a larger value of g_2 than envisaged by the general bag model dependent analysis of Lie-Svendson and Høgaasen [62], one has to examine any possibly misleading concepts of the 'bag'. A comparison of the g_A/g_V ratios in the equations (3.21) with (3.24-29) shows that reasonably good agreement can be found at least in the case of the chiral bag model dependent calculations. The importance of including chiral symmetry into the bag model can also be seen by some chiral bag model explanations of a reduced quark spin sum [67].

The possible future confirmation of non-vanishing induced second-class form factors and of the consequent shifts of the g_A/g_V ratios in the semileptonic hyperon decays is also of great importance for the judgement of Lipkin's claim that SU(3) flavour symmetry deserves 'a decent burial and a honorable place in history' [68].
Lipkin's disappointment with SU(3) flavour stems mainly from the following argument [69].
The surprising agreement of $\frac{g_A}{g_V}|_{\Sigma^-\to n} = -0.328 \pm 0.019$ with the NQM prediction $\frac{g_A}{g_V}|_{\Sigma^-\to n} = -\frac{1}{3}$ contradicts the strong reduction in case of the neutron β-decay from its NQM value $\frac{g_A}{g_V}|_{n\to p} = \frac{5}{3}$ to its experimental value $\frac{g_A}{g_V}|_{n\to p} = \frac{5}{4}$. Lipkin explains this discrepancy by introducing very different SU(6) breaking admixtures into the wavefunctions. He argues that the transitions are all very different. The transition $\Sigma^- \to n$ should be unique in its simplicity, because it allows the clear separation of one active part (the $s \to u$ transition) from an inert spectator part. Lipkin decomposes the wavefunction of the Σ^- and the neutron with the mixing parameters θ^Σ and θ^n as follows:

$$\begin{aligned} |\Sigma> &= cos\Theta^\Sigma |S_1^\Sigma; s> + sin\Theta^\Sigma |S_0^\Sigma; s> \\ |n> &= cos\Theta^n |S_1^n; u> + sin\Theta^n |S_0^n; u> \end{aligned} \tag{3.30}$$

where $S_0^\Sigma, S_1^\Sigma, S_0^n, S_1^n$ denote the spectator states in the Σ and the neutron, with angular momentum zero and one respectively.
According to Lipkin this simplicity of the Σ decay should be seen in contrast to the neutron decay, where there are two active $d \to u$ transitions and consequently more free parameters. But when SU(6) is broken in such a radical manner, it obviously does not make sense to keep SU(3) symmetry. For this reason Lipkin calls the

SU(3) characteristic F/D ratio a 'fudge factor' [68]. But taking into account the new data of $\frac{g_A}{g_V}|_{\Sigma^- \to n}$ his argument has to be revised.

It is a hopeful idea that with the new z or F/D value, which is very close to its SU(6) symmetric value, hadrons can still be SU(3) eigenstates without any mixing of different SU(3) multiplets. SU(3) flavour symmetry is crucial for our understanding of spectroscopy and before it is given up completely, a minimal SU(3) breaking pattern should be studied which preserves the most essential properties of this symmetry.

Another intuitive argument against exact SU(3) symmetry in the polarisation (and consequently in the baryon wavefunctions) is Lipkin's suspicion concerning the small strange quark polarisation in the nucleon in comparison with e.g. the up quark polarisation in the Σ^- [69]. This would imply the failure of the ansatz in equation (3.22). The equivalent SU(3) breaking of the sea quark polarisation inside the nucleon

$$\triangle u_{sea} = \triangle d_{sea} = (1+\epsilon)\triangle s, \text{ with } \epsilon > 0 \tag{3.31}$$

can however only be proven if, for some peculiar reason, the production of strange quark pairs inside the nucleon remains suppressed completely, i.e. $< p|\bar{s}s|p >= 0$. This is a very unlikely situation, as is made clear by many experiments which have indicated the possibility of strangeness inside the nucleon (see also chapter 5). The basic picture we have of sea quark production is given by the arguments which led to perturbative QCD. The flavour blind splitting functions P_{qG}, P_{Gq} do not depend on the quark masses, which are kept outside the perturbative treatment. The vector coupling γ_μ of the gluon does not allow a helicity flip of any produced $q\bar{q}$ pair, being SU(3) invariant. The gluon spin is taken up by the angular momentum $L_{q\bar{q}}$ of the

$q\bar{q}$ pair [70], and an (anti)quark spin flip might occur non-perturbatively (with a probability $\propto m_q^2$) to compensate [71]. It is therefore difficult to follow Lipkin's argument [68,69] as to why, at current quark level, the sea quark polarisation should be only SU(2) invariant with the strange quark component being suppressed.

With regard to the EMC experiment there is still the dilemma of dynamically explaining a negative strange quark polarisation in the presence of a mostly positive gluon polarisation. Before this problem is taken up again, another aspect of SU(3) symmetry breaking should be discussed which is also of relevance for the evaluation of the EMC result.

In his criticism of the SU(3) flavour symmetry Lipkin pointed out [21],[68],[69] that the SU(3) relation between the charged axial current operator and the neutral axial current operator

$$< n|\bar{u}(1+\gamma_5)\gamma_\mu s|\Sigma^- > + < n|\bar{s}(1+\gamma_5)\gamma_\mu u|\Sigma^- >= \triangle d - \triangle s \tag{3.32}$$

is violated as soon as a neutral sea violates SU(3) with $\epsilon > 0$ in (3.31). This happens even if the matrix elements of the Cabibbo charged axial current operator between different states in the baryon octet remain related by SU(3). It is of course questionable how large such an impurity could be, and it seems reasonable to take this effect to be only of second order. It might be of interest, that despite his criticism against the use of SU(3) in evaluating a_0 from the EMC experiment, Lipkin arrives more or less at the same value for a_0, obtained with SU(3) symmetry intact.

There is obviously a clear advantage in keeping SU(3) symmetry unviolated as far as baryon wavefunctions and axial vector currents are concerned. No essential predictive power has to be given up and the weak currents remain independent of the quark masses.

The current quark masses ($m_q \neq 0$) and the chiral condensate ($< q\bar{q} > \neq 0$) break the chiral $SU(3) \times SU(3)$ symmetry and may generate the constituent quarks (see chapter 5); but the way SU(3) is broken explicitly by a different strange quark mass $m_u = m_d \neq m_s$ at the constituent quark level, which gives rise to an 'induced second-class' form factor, does not imply the existence of weak 'second-class currents'. The standard left-handed quark current, which is a 'first-class' current, is the only weak current needed to describe the weak hadronic interaction in the framework of the standard model.

Finally, another aspect should be discussed which will bring us back to the problem of the quark spin sum. In equations (3.26/7) a new value of F/D or equivalently z was derived. There has been for some time controversy over the validity of QCD, which culminated in Close's defence of QCD [19] where he argued that Preparata and Soffer [72] used the wrong value of F/D. The value used was $F/D = 2/3$ which is in agreement with the value in equation (3.27). Close's value of $F/D = 0.56$ is taken from an earlier analysis of Close and Roberts [73] where they argued for the single reference value $\frac{g_A}{g_V}|_{\Sigma^- \to n} = -0.328 \pm 0.019$ neglecting all other measured hyperon decays, because the value from the Σ^--decay was the most precisely measured. But this is true only if one sets $g_2 = 0$. Does the new value of F/D now jeopardise QCD? Definitely not!

Ellis and Karliner [74] have already argued that the inequality

$$|\triangle s(x)| \leq s(x) \tag{3.33}$$

does not give a suitable constraint on $\triangle s$ due to diffraction: $s(x) \sim x^{-1}$ for $x \to 0$. This is the reason why the DIS neutrino data, measuring the strangeness inside the proton, cannot provide a stringent bound per se. Instead one has to construct a non-diffractive distribution function s_{ND} with the infinity removed. A prescription of how to do this was given by Ellis and Karliner [74]. Because of the existence of the gluon anomaly, QCD has never been put in doubt by the EMC experiment. When Close [19] discusses the dependence of $\triangle s$ on F/D with the EMC result as input

$$\triangle s \simeq -(F/D - 0.40) \pm 0.07 \tag{3.34}$$

he neglects the QCD-corrected value of $\triangle s - \frac{\alpha_s}{2\pi}\triangle G$, which can already be enough to balance any disagreement between polarised and unpolarised distribution functions.

Close's claim that the EMC analysis is very sensitive to the F/D value is not quite plausible if one regards the range of z, which is in favourable agreement with a vanishing of the flavor singlet axial-vector-current matrix-element.

$$\int_0^1 g_1^p dx = \frac{1}{12}(1 - \frac{\alpha_s}{\pi})(1 + z_{DIS})a_3 \tag{3.35}$$

Taking into account the radiative corrections with the value $\alpha_s(Q^2) = 0.25 \pm 0.02$ and using the combined EMC and SLAC data for $\int_0^1 g_1^p dx$ [4], one finds

$$z_{DIS} = 0.303 \pm 0.104 \pm 0.155 \tag{3.36}$$

where the errors are statistical and systematic respectively. So, the experiment does not appear to be especially sensitive on a variation of z or F/D.

The spin problem remains a question of how to generate a negative strange quark polarisation inside the nucleon, however small. The apparent difference between current and constituent quarks remains a problem. Perturbative QCD might not be sufficient for its solution, but its validity in this respect is beyond any doubt; not least because of the discussion in the next chapter, which will show the existence of negative strange quark polarisation already present at the constituent quark level. In addition to this surprising observation, the new value of z or F/D (which implies unbroken SU(3) symmetry in the baryon wavefunctions and in the axial vector currents) will be confirmed in the next chapter in our analysis of the baryon magnetic moments.

4. The magnetic moments of the baryons

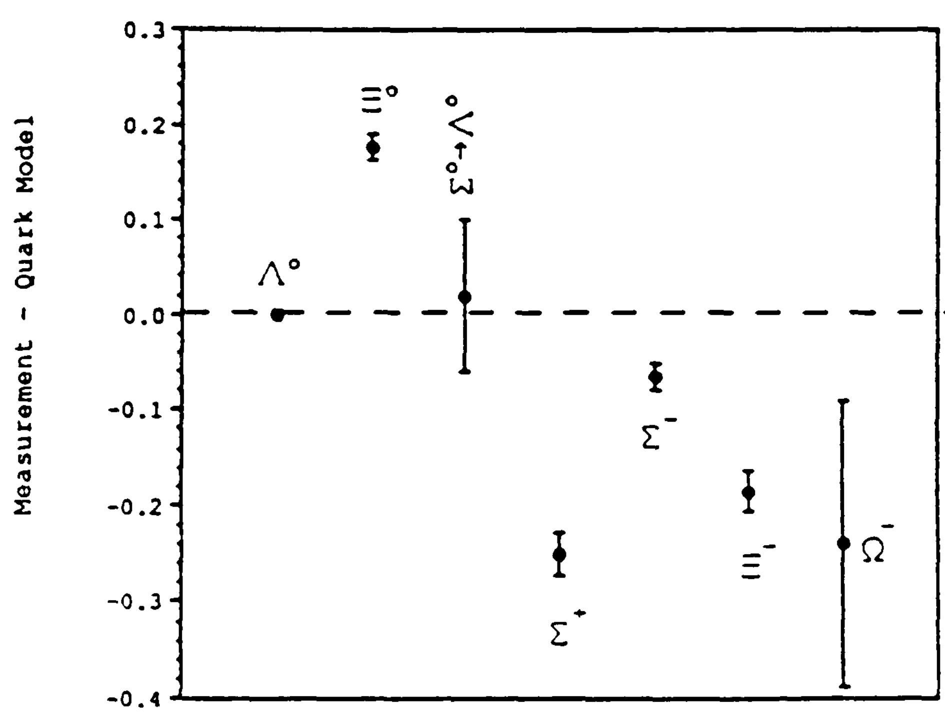
4.1. The constituent quark spin contribution

Often cited as the experimental evidence for the NQM and the SU(6) symmetry, the data [7][75] as well as the theoretical information [20,68] [76–83] on magnetic moments of baryons have been reviewed in the light of the EMC result [4]. In ref.[7] a striking deviation of as much as 37.76% between the naive quark model prediction and the measurement of the magnetic moment of the Ξ^- has been pointed out.

This discrepancy between the experimental data and the NQM cannot be understood even after several systematic corrections are made, e.g.:
Isgur and Karl [84] studied the influence of possible configuration mixing effects between the 56-plet and the 70-plet on baryon magnetic moments; they included relativistic bag model corrections by assuming that the quarks are confined in cavity under outside pressure, and they even took isospin violations into account. But because these additional considerations could not give a satisfactory improvement between theory and experiment, Isgur and Karl had to conclude that the details of this disagreement depend on too many unknowns.
It may be worth pointing out that the above corrections have always been believed to be of a perturbative nature, because the NQM was still seen as a perfect first order approximation to the data.

However, the precision of the measurements has improved considerably in recent years, and the gap between experiment and NQM based theory has widened [7]:

	Magnetic Moment μN		Quark Model μN	Difference μN		σ	%Dif
p	2.792847386	± 0.000000063	input				
n	-1.91304275	± 0.00000045	input				
Λ°	-0.613	± 0.004	input				
Σ^{+}	2.419	± 0.022	2.67	-0.251	± 0.022	-11.41	-9.40
Σ^{-}	-1.156	± 0.014	-1.09	-0.066	± 0.014	-4.71	6.06
$\Sigma^{\circ}\rightarrow\Lambda^{\circ}$	-1.61	± 0.08	-1.63	0.02	± 0.08	0.25	-1.23
Ξ°	-1.253	± 0.014	-1.43	0.177	± 0.014	12.64	-12.38
Ξ^{-}	-0.675	± 0.022	-0.49	-0.185	± 0.022	-8.41	37.76
Ω^{-}	-2.08	± 0.15	-1.84	-0.24	± 0.15	-1.60	13.04



In the light of the EMC experiment it was therefore natural to ask if the magnetic moments of the baryons really do contradict the result of a vanishing quark spin sum. The gluon anomaly might play a more significant role at the current quark level than at the constituent quark level, but the QCD perturbatively calculated evolution of $a_0(Q^2)$ only suggests a change of order α_s (see equation 2.34). Strictly speaking, the extrapolation of this perturbatively calculated evolution behaviour has to be made into the non-perturbative region. This is known to work well in the unpolarised case.

The subsequent discussion of Karl's most recent analysis of the magnetic moments [76] will be compared with the discussion of my own work [23], because our conclusions are basically the same. Added weight may be given to the common conclusion since it was arrived at by quite different means.
The starting point for Karl's discussion is the specification of the class of models which allow one to connect the matrix elements of the axial vector current operator to the magnetic moment operators. In a class of shell models which include the NQM, the bag model and their simple extensions which contain antiquarks as well, Karl defines one-body magnetic moment operators μ_q^* for every quark and antiquark, the occupation number of u quarks with $J_z = +\frac{1}{2}$ $n(u_\uparrow)$, and the occupation number of $\bar{u}$ quarks with $J_z = +\frac{1}{2}$ $n(\bar{u}_\uparrow)$, and so on. All quarks and antiquarks of a given flavour are assumed to be in a singular mode of angular momentum one-half of some potential, or spherical cavity. The sum of all these one-body operators is the magnetic moment operator of the baryon $M_z(B)$, so that in the case of the proton

$$M_z(p) = M_z(p;u) + M_z(p;d) + M_z(p;s) \tag{4.1}$$

where

$$M_z(p;q) = [n(q^{\uparrow}) - n(q^{\downarrow}) - n(\bar{q}^{\uparrow}) + n(\bar{q}^{\downarrow})]\mu_q^* \tag{4.2}$$

Similarly he introduces one-body matrix elements a_q of the axial vector current for every quark and antiquark, so that

$$\triangle q = [n(q^{\uparrow}) - n(q^{\downarrow}) + n(\bar{q}^{\uparrow}) - n(\bar{q}^{\downarrow})]a_q \tag{4.3}$$

Weinberg [18] has given an even more general justification for the sum of the quark magnetic moment operators being the baryon magnetic moment operator. This explanation concerns Dirac nature of a constituent quark. Starting from an effective chiral Lagrangian (see next chapter) Weinberg proves in the limit of large N_c, which is an approximation to the $N_c = 3$ case of QCD, that the constituent quark has no anomalous magnetic moment, $\kappa_q = 0$, and that the weak couplings of a constituent quark and of a current quark are the same (i.e. $g_A = 1$). His argument, which gives the constituent quark magnetic moment its Dirac value

$$\mu_q = \frac{e_q}{2m_q} \tag{4.4}$$

and the same g_A of a current quark does not rely on the renormalisability of the theory which rules out a Pauli magnetic moment in the QED Lagrangian.
The corrections to $g_A = 1$ and to $\kappa_q = 0$ in $\frac{1}{N_c}$ have been shown [85] to be small and in the case of κ_q even negligible.
The authors of ref.[18,85] argued that constituent quarks acquire their masses from the spontaneous chiral symmetry breaking. In order to treat these quarks as massive

particles pions had to be introduced among the degrees of freedom of the constituent quark model [86]. For the calculation of g_A for a quark, the analogue of the Adler-Weisberger sum-rule [87] for quark-pion scattering

$$g_A^2 = 1 - \frac{2f_\pi^2}{\pi}\int_0^\infty \frac{d\omega}{\omega}[\sigma_-(\omega) - \sigma_+(\omega)] \tag{4.5}$$

has been used, where $\sigma_\pm(\omega)$ is the total cross-section for the scattering of a $\pi^\pm$ of energy ω on a quark at rest. For the calculation of κ_q, the analogue of the Drell-Hearn-Gerasimov sum rule [88]

$$\kappa_q^2 = \frac{m^2}{2\pi^2\alpha_{em}z^2}\int_t^\infty \frac{d\omega}{\omega}[\sigma_P(\omega) - \sigma_A(\omega)] \tag{4.6}$$

for photon-quark scattering has been used, where σ_P and σ_A represent the cross-sections for a parallel and antiparallel photon and quark spins, and ω denotes the incoming energy of the photon in the frame where the target quark is at rest. ze denotes the charge of the target and t is the threshold for the relevant process. As it turns out from the analysis [18,85], both integrals in (4.5) and (4.6) do more or less vanish.

However, in the case that a constituent quark does have these elementary properties of a bare Dirac particle (especially (4.4)), one may legitimately ask why the baryonic data do not coincide better with the predictions.

In generalising Weinberg's argument from the two flavour case to the three flavour case, we may use the fact that in the chiral Lagrangian density the massive constituent quarks and the pseudoscalar mesons (treated as NG bosons) are subject to the colour force only at large separations (see Manohar and Georgi [86]). The following set of equations should then hold to good accuracy:

$$\mu_B = \sum_{q=u,d,s} e_q \frac{\sigma_q}{2m_q} \tag{4.7}$$

where m_q is the effective quark mass and $\sigma_q = q^\uparrow - q^\downarrow - (\bar{q}^\uparrow - \bar{q}^\downarrow)$ is the relevant combination of the fractions of the proton spin carried by the constituent quark q and anti-quark $\bar{q}$ for the magnetic moments. This should not be confused with the combination $\triangle q =< p^\uparrow|\bar{q}\gamma_\mu\gamma_5 q|p^\uparrow >= q^\uparrow - q^\downarrow + (\bar{q}^\uparrow - \bar{q}^\downarrow)$ measured in the polarized deep inelastic scattering experiments, where the photon current-quark cross-section is sensitive to the polarization of the quark but not to the sign of the electric charge. Because of the different charge conjugation of the magnetic moments and the axial vector currents it is not possible to replace σ_q with $\triangle q$, even in the case that $g_A = 1$. Under the assumption of unbroken SU(3) flavour symmetry in the baryon wavefunction, it is however possible to relate the flavour octet matrix elements:

$$\begin{gathered} \sigma_u - \sigma_d \sim \triangle u - \triangle d = a_3 \\ \frac{1}{\sqrt{3}}(\sigma_u + \sigma_d - 2\sigma_s) \sim \frac{1}{\sqrt{3}}(\triangle u + \triangle d - 2\triangle s) = a_8 \end{gathered} \tag{4.8}$$

For this reason z or the F/D-ratio should stay the same, if calculated from the β-decays or from the magnetic moments of the baryons. For constituent quarks being bare Dirac particles with $g_A = 1$ these relations reduce to identities. But one has to be careful about the explanation of the reduced value of a_3 from its NQM value 5/3 to its experimental value 5/4.

Normally this reduction is attributed to the 'lower' component of the relativistically boosted quark spinor along the z-axis, as discussed in the Introduction. In this frame, however, additional antiquarks added to the normal three quark baryon display a different transformation behaviour due to the intrinsic parity operator β

being multiplied with the σ_x and σ_y generators of SU(2). The resulting $SU(2)_W$ group commutes with the Lorentz-boost α_3. In the corresponding $SU(6)_W$ symmetry of the quark currents, mesons are classified in a different way than they are in the NQM. Due to the antiquarks the classification of the vector mesons and pseudoscalar with respect to the spin is interchanged. This is also the reason underlying the explanation of the non-leptonic decays of the $\frac{3}{2}^+$ decuplet and 1^- vector meson states (1.21), which are forbidden in the NQM (1.18). These decays are still difficult to understand in the framework of chiral Lagrangians, because in the limit of massless NG bosons (i.e. massless pions) such a transformation is not possible.
The link between the one-body axial-vector current and the magnetic moment operators made by Karl becomes questionable as soon as the contribution of the polarised antiquarks are parametrised by a constant ratio λ_q:

$$n(\bar{q}^{\uparrow}) - n(\bar{q}^{\downarrow}) = \lambda_q[n(q^{\uparrow}) - n(q^{\downarrow})] \tag{4.9}$$

This danger was also noticed by Jaffe and Lipkin [89] who had to exempt from taking part in any β-decay, the additional $Q\bar{Q}$ pair which they introduced recently into a NQM baryon.

Based on Weinberg's argument that $\kappa_q = 0$ for the constituent quarks, it is reasonable to first evaluate the following equations only with respect to σ_q without any consequence for $\triangle q$:

$$
\begin{aligned}
\mu_p &= \mu_u\sigma_u + \mu_d\sigma_d + \mu_s\sigma_s \\
\mu_n &= \mu_d\sigma_u + \mu_u\sigma_d + \mu_s\sigma_s \\
\mu_{\Sigma^+} &= \mu_u\sigma_u + \mu_s\sigma_d + \mu_d\sigma_s \\
\mu_{\Sigma^-} &= \mu_d\sigma_u + \mu_s\sigma_d + \mu_u\sigma_s \\
\mu_{\Xi^\circ} &= \mu_s\sigma_u + \mu_u\sigma_d + \mu_d\sigma_s \\
\mu_{\Xi^-} &= \mu_s\sigma_u + \mu_d\sigma_d + \mu_u\sigma_s \\
\mu_\Lambda &= \frac{1}{6}(\sigma_u + 4\sigma_d + \sigma_s)(\mu_u + \mu_d) + \frac{1}{6}(4\sigma_u - 2\sigma_d + 4\sigma_s)\mu_s \\
\mu_{\Sigma\Lambda} &= \frac{1}{2\sqrt{3}}(\sigma_u - 2\sigma_d + \sigma_s)(\mu_u - \mu_d)
\end{aligned}
\tag{4.10}
$$

The form of these equations, in which the baryons are assumed to be SU(3) eigenstates, is of course the same in Karl's analysis and, following his observation, resemble the equations for the baryon octet masses in chiral perturbation theory. Analogous sum-rules can be written down:

$$
\begin{aligned}
&\mu_{\Sigma^+} + \mu_\Lambda + \mu_{\Sigma^-} = 2\mu_p + 2\mu_n + 2\mu_{\Xi^\circ} + 2\mu_{\Xi^-} \\
&[-2.34 \pm 0.06\mu_N \neq -2.10 \pm 0.04\mu_N]
\end{aligned}
\tag{4.11}
$$

and

$$
\begin{aligned}
&\mu_n - \mu_p + \mu_{\Sigma^+} - \mu_{\Sigma^-} - \mu_{\Xi^\circ} + \mu_{\Xi^-} = 0 \\
&[-0.49 \pm 0.06\mu_N \neq 0]
\end{aligned}
\tag{4.12}
$$

As indicated in square brackets the experimental data do not support these sum-rules. This shows that there are no exact solutions to this set of equations as it stands. Because this set cannot fit the data with an accuracy approaching the experimental error, Karl introduces in his article a 'theoretical error' which should express how close one might come to the experimental measurements. From the failure of the sum-rule this 'theoretical error' is estimated to be $0.1\mu_N$ and is added to the experimental error of each data in quadrature.

Because of the symmetric form of equations (4.10), the real number of free parameters is reduced to four independent quantities which Karl chooses to be (recalling his identification $\triangle q \sim \sigma_q$)

$$\triangle u + \triangle d + \triangle s \ , \quad \frac{1}{\sqrt{3}}(\triangle u + \triangle d - 2\triangle s) \ , \quad \mu_d \ , \quad \mu_s \tag{4.13}$$

In the four-dimensional space of these parameters Karl searches for the minimum of the quantity χ^2:

$$\chi^2 = \sum_{i=1}^{8} \frac{(Eq.(4.10) - expt.)_i^2}{(error)_i^2} \tag{4.14}$$

Without the introduction of the 'theoretical error', the fits are better for the best measured magnetic moments p, n, Λ at the expense of those of other baryons, but the result of the analysis does not change significantly. Probably because of relativistic corrections, Karl believes that the set of equations (4.10) should actually fit better the Ξ's with more heavy quarks.
The main result of Karl's numerical analysis is that

$$\triangle u + \triangle d + \triangle s = 0.27 \pm 0.23 \tag{4.15}$$

is more compatible with the EMC experiment than with the NQM, even if the absolute quark spin sum is determined with little accuracy.

The set of equations (4.10) has been studied by several authors [80][81][82][83]who changed all kinds of parameters, but no analysis could explain the systematic 'theoretical error' which Karl had to introduce. Another deficiency in the standard analysis of these equations is that a 'fit' is either over-parametrised, or leaves data unexplained (mostly $\mu_{\Sigma^\pm}$).

In the following, a systematic study is presented [23], in which the SU(3) flavour symmetry remains unbroken as much as possible, and only the most plausible breaking mechanism (i.e. via the quark masses) introduces a new set of parameters, without violating SU(2). In this way a prescription can be given, which keeps the number of free parameters to a minimum. This work indeed pre-dates Karl's analysis, but his result, which is basically the same, supports the pattern in which we observe the symmetry breaking to take place.

In the previous chapter we discussed what effect a non-vanishing g_2 term has in semileptonic hyperon decays and how it could be interpreted. The most convincing interpretation, due to its simplicity, is that SU(3) flavour symmetry is broken by $m_s \neq m_d = m_u$ without violating SU(3) symmetry in the baryon wavefunction. Taking into account the measured value of g_2 in the $\Sigma^- \to n$ decay and the consequently modified value of g_A/g_V, one finds that the new value of F/D is again in agreement with its SU(6) value. This is the reason why, even with respect to Lipkin's objection [21] to the use of SU(3) while breaking SU(6), there is now no more reason to break SU(3) in the wavefunction. No essential predictive power has to be given up.

In the same sense, of course, the SU(3) symmetry in the baryon wavefunction should also not be affected by unequal quark masses when the magnetic moment operators are applied.

The SU(3) breaking is parametrised by the constant factor S:

$$m_u = m_d = \mathcal{S} \cdot m_s \ , \qquad 0 < \mathcal{S} \leq 1 \tag{4.16}$$

The isospin symmetry should not be violated by any first order analysis. This is also the reason why the set of equations (4.10) should be evaluated in an SU(2) invariant way!

In order to concentrate on the symmetry of the wavefunctions, first the a *priori* unknown quark masses have to be eliminated. Because Iso-spin is a better symmetry compared to U-spin or V-spin, it is better to study states with different strangeness quantum number. In fact, Lipkin [21] stresses the need to distinguish between relations which require only isospin and those requiring the full SU(3).

Because the magnetic moments of the isoscalar Λ and the magnetic transition $\Sigma^o - \Lambda$ are not related in their SU(3) symmetric Clebsch-Gordan coefficient decomposition by a simple isospin transformation to any other magnetic moment, we postpone a discussion of them until later.

The solution of every pair of equations is readily obtained by looking in an isospin invariant manner at the ratios:

$$n_1 = \frac{\mu_p + \mu_n}{\mu_p - \mu_n} = 0.187$$

$$n_2 = \frac{\mu_{\Sigma^+} + \mu_{\Sigma^-}}{\mu_{\Sigma^+} - \mu_{\Sigma^-}} = 0.353 \pm 0.007 \tag{4.17}$$

$$n_3 = \frac{\mu_{\Xi^o} + \mu_{\Xi^-}}{\mu_{\Xi^o} - \mu_{\Xi^-}} = 3.159 \pm 0.053$$

The three experimental numbers $n_1 = 0.187$, $n_2 = 0.353 \pm 0.007$, $n_3 = 3.159 \pm 0.053$ are related to the σ_q's in the following way:

$$n_1 = \frac{\sigma_u + \sigma_d - 2S\sigma_s}{3(\sigma_u - \sigma_d)} = 0.187$$

$$n_2 = \frac{\sigma_u + \sigma_s - 2S\sigma_d}{3(\sigma_u - \sigma_s)} = 0.353 \pm 0.007 \tag{4.18}$$

$$n_3 = \frac{\sigma_d + \sigma_s - 2S\sigma_u}{3(\sigma_d - \sigma_s)} = 3.159 \pm 0.053$$

This set of three equations can be solved for the two free parameters $S = \frac{m_u}{m_s}$ and z_{mag}

$$z_{mag} = \frac{\sigma_u + \sigma_d - 2\sigma_s}{3(\sigma_u - \sigma_d)} \tag{4.19}$$

The result is

$$z_{mag} = \frac{(n_3 + n_2 - 2\sqrt{n_1 n_2 + n_2 n_3 - n_1 n_3})}{3(n_3 - n_2)} = 0.235 \pm 0.002 \tag{4.20}$$

which would correspond to

$$F/D = 0.743 \pm 0.005 \tag{4.21}$$

a value in clear agreement with Karl's result $z = 0.23$, and with (3.26/27). Karl's subsequent prediction for $\frac{m_d}{m_s}$, however, just "remains rather close to 0.6, the value traditionally found in quark model fits, which equals the constituent quark mass ratio".

Our result for S is more precise:

$$\mathcal{S} = \frac{3}{2}\sqrt{n_1 n_2 + n_2 n_3 - n_1 n_3} - \frac{1}{2} = 0.653 \pm 0.024 \tag{4.22}$$

A result which is by no means trivial since the general nature of the ansatz did not include any mass scale!

Strictly speaking, only eq.(4.22) allows the interpretation of $\mathcal{S}$ as the constituent quark mass ratio $m_u : m_s$.

It is clear that the value of this ratio, in addition to the equality of m_u and m_d, is precisely what is expected from a constituent quark. By analogy with the current quarks which are most accurately described by their mass ratios calculated from the chiral perturbative analysis [90,91], we do not need the absolute values of the quark masses in order to decide whether constituent or current quarks play the fundamental role in observed processes. The characteristic mass ratio in eq.(4.22) is a clear signal for constituent quarks and is consistent with the baryon mass spectrum. $\mathcal{S} = 0.622 = \frac{2(m_{\Sigma^*} - m_\Sigma)}{2m_{\Sigma^*} + m_\Sigma - 3m_\Lambda}$ is, for example, the result of a 'QCD' calculation, performed for s-wave baryons [92]. Also in dynamical processes, constituent quarks appear as massive objects recognizable through this particular mass ratio. $\mathcal{S} = 0.70 \pm 0.07$ is, for example, the result of an analysis of the radiative decays of charged and neutral K^* vector mesons [93].

The most surprising result, however, is the quark spin sum for which we find

$$\sigma_u + \sigma_d + \sigma_s = (0.086 \pm 0.028)(\sigma_u - \sigma_d) \tag{4.23}$$

and for each flavor separately:

$$
\begin{aligned}
u^{\uparrow} - u^{\downarrow} &= +(0.646 \pm 0.010)(\sigma_u - \sigma_d) + \bar{u}^{\uparrow} - \bar{u}^{\downarrow} \\
d^{\uparrow} - d^{\downarrow} &= -(0.354 \pm 0.010)(\sigma_u - \sigma_d) + \bar{d}^{\uparrow} - \bar{d}^{\downarrow} \\
s^{\uparrow} - s^{\downarrow} &= -(0.206 \pm 0.008)(\sigma_u - \sigma_d) + \bar{s}^{\uparrow} - \bar{s}^{\downarrow}
\end{aligned}
\tag{4.24}
$$

This is indeed a surprising observation. A non-vanishing expectation value of polarized strange or anti-strange quarks at the constituent quark level in a proton is contributing to its magnetic moment!

No anomalous gluon contribution can change this result. In this respect, the flavour singlet quark spin sum has no ambiguity in its definition at the constituent quark level.

A priori, there are no constraints for $\sigma_u - \sigma_d$, besides being positive. In the NQM the value of $\sigma_u - \sigma_d$ is 5/3. The experimental uncertainty in $\sigma_u - \sigma_d$ will probably remain for the time being. The freedom however to set $\sigma_u - \sigma_d = a_3$ in the set of equations (4.10) without changing the analysis can be considered a success. No more 'small' quark spinor component is needed for the 25% reduction of $\sigma_u - \sigma_d$ from its NQM value to a_3, measured in the β-decay of the neutron! Weinberg could not explain this reduction on more general grounds [18].
The problem with the different transformation behaviour of any additional anti-quarks, and their contrived exclusion from any active role in the weak interaction, might be solved. By linking directly $\sigma_u - \sigma_d$ to a_3, Karl implicitly anticipated this freedom. On the other hand, Karl could not explain the large value of $z = 0.23$ which should be the same value for the hyperon decays, especially in the class of shell models he considers. He did not know [94] of the new measurement of $\frac{g_1}{f_1}$ in the Σ^- decay. Only with this additional input, is Karl's analysis consistent.

4.2. The environment dependent constituent quark mass

Why could this system of relations (4.10) so easily be satisfied, despite the failure of the second sum-rule (4.12) which involves the same data?

Our ignorance is mainly about the exact values of the quark masses, and by solving the set of equations (4.10) in an isospin invariant manner, one strangeness or baryon mass dependent scaling factor has been dropped. Already in an article of 1980, Isgur and Karl [84] wrote that 'there is no reason to suppose that the magnetic moment of a given quark bound in one baryon will be identical to the moment of the same quark in a different baryon'. By evaluating the equations in an SU(2) invariant way we have allowed the quark masses to scale. The SU(3) symmetry breaking factor $S = 0.653 \pm 0.024$ remains constant, but the absolute values of all constituent quark masses depend on the environment!

There is no clear understanding how the constituent quarks are generated in terms of the fundamental theory of strong interaction, QCD. In spite of a rather vague description of a constituent quark, it is however usually assumed that the masses of these complex entities remain constant, independent of their environment, not even allowing a mass fluctuation of a few MeV.

This seems to be unnatural and also leads in addition to discrepancies between the experimental data and the theoretical sum rules (4.11) and (4.12) involving magnetic moments.

As a consequence of this, one has not had so far any reliable derivation of the quark spin expectation values contributing to the magnetic moments. The small experimental errors give constraints which rule out existing models and impose a 'theoretical error' on our ignorance.

A possible dynamical mechanism of the SSB of chiral symmetry is given by the Nambu Jona-Lasinio (NJL) model [95] which will be discussed in further detail in the next chapter. In the framework of this model it is possible to calculate the dynamical constituent quark masses. If the gap equations which define these masses (see next chapter) describe particles in different hadronic environments, it seems natural to expect different quark masses without losing any symmetry in the wavefunction or in the interaction.

It is encouraging to observe the behaviour of the resulting scaling factor of the quark masses:

$$f_\Sigma := \frac{m_u|_\Sigma}{m_u|_N} = 1.117 \pm 0.009 \ , \quad f_\Xi := \frac{m_u|_\Xi}{m_u|_N} = 1.150 \pm 0.033 \tag{4.25}$$

This behaviour is exactly reflected by the hadron masses

$$m_\Sigma \simeq 106\% \frac{m_N + m_\Xi}{2} \tag{4.26}$$

and can be seen as a strong support for the above argument. Instead of the above sum-rules (4.11) and (4.12), an interesting quadratic relation emerges from (4.10):

$$\mu_p^2 - \mu_n^2 + f_\Xi^2(\mu_{\Xi^o}^2 - \mu_{\Xi^-}^2) = f_\Sigma^2(\mu_{\Sigma^+}^2 - \mu_{\Sigma^-}^2) \tag{4.27}$$

At this stage it is necessary to include the other measured magnetic moments of the Λ and the $\Lambda\Sigma$-transition into the argument. The magnetic moment of the Σ^o has not yet been measured but it is straightforwardly predicted to be

$$\mu_{\Sigma^o} = \frac{1}{2}(\mu_{\Sigma^+} + \mu_{\Sigma^-}) \tag{4.28}$$

In contrast with this unmeasured moment, the magnetic moment of the Λ has been measured to great accuracy. In its SU(3) decomposition (with the measured moment in boldface)

$$(\mathbf{-0.613 \pm 0.004}) \cdot 3 \cdot \frac{m_u}{m_p}\Big|_{\Lambda} = \frac{1}{6}(4\sigma_d + \sigma_u + \sigma_s - \mathcal{S}(4\sigma_u + 4\sigma_s - 2\sigma_d)) \tag{4.29}$$

the uncertainties of the experimental error and the uncertainties about the spin expectation values σ_q and S, however, add up if we evaluate the above relation with respect to the mass scaling factor f_Λ:

$$f_\Lambda := \frac{m_u|_\Lambda}{m_u|_N} = 1.028 \pm 0.056 \tag{4.30}$$

Nevertheless, this number is again in perfect agreement with the observation that the mass of the Λ is between the mass of the nucleon and the mass of the Σ, lending further support to the view that the notion of an environment dependent constituent quark is meaningful.

The magnetic transition moment $\mu_{\Lambda\Sigma}$ appears to be independent of SU(3) symmetry breaking effects: The S parameter does not enter the SU(3) symmetric decomposition

$$(\mathbf{1.61 \pm 0.08}) \cdot 3\frac{m_u}{m_p}\Big|_{\Lambda\Sigma} = \frac{1}{2\sqrt{3}}(2\sigma_d - \sigma_u - \sigma_s) \tag{4.31}$$

The corresponding effective quark mass fits rather in a Λ than in a Σ^o environment:

$$f_{\Lambda\Sigma} := \frac{m_u|_{\Lambda\Sigma}}{m_u|_N} = 0.97 \pm 0.05 \tag{4.32}$$

The difference between the Λ and the Σ mass can be understood as a consequence of a QCD induced spin-spin hyperfine interaction. Being inversely proportional to the product of the effective quark masses of the interacting quarks and otherwise flavour independent, the spin dependent hyperfine interaction

$$v_{ij}^{hyp} = \frac{\vec{\sigma}_i \cdot \vec{\sigma}_j}{m_i m_j} \bar{v}_{ij}^o \tag{4.33}$$

was introduced on a purely phenomenological basis in order to better understand hadron masses and their magnetic moments. The conjecture [96] that the hadron and effective quark masses are split into two terms

$$m_B = \sum_i m_i + \sum_{i>j} v_{ij}^{hyp} \tag{4.34}$$

$$\sum_i m_i = \sum_i \epsilon_i + \sum_{i>j} v_{ij}^o \tag{4.35}$$

where ϵ_i is the full single particle energy and v_{ij}^o is a flavour independent part of the two body interaction, could be simply modified by taking the flavour dependent part $\sum_{i \neq j} v_{ij}^{hyp}$ into the definition of the effective quark mass:

$$m_i = \epsilon_i + \sum_{i \neq j} v_{ij}^o + \sum_{i \neq j} v_{ij}^{hyp} \tag{4.36}$$

After switching off spin-spin interactions all constituent quark masses rescale back to their 'original' groundstate value. A rough estimate of their value can be obtained simply by reducing the unregular baryon mass gap in the flavour octet through $\frac{m_B}{f_B}$ to approximately equidistant mass gaps

$$(m_s - m_u)|_p = (103 \pm 37)MeV \tag{4.37}$$

Taking $S = \frac{2}{3}$ one had $m_u|_p = (206 \pm 74)MeV$ which is to be compared with the absolute constituent quark mass scale in the nucleon:

$$m_u|_N \simeq (\sigma_u - \sigma_d) \times 200MeV \tag{4.38}$$

The above analysis raises the question as to why the effective quark mass should include the hyperfine interaction (4.36). An answer to this problem rests ultimately in dynamical mass generation, but the following phenomenological consideration should be noted:

The hyperfine interaction (4.33) is seen as the reason for the SU(6) breaking of the spin-flavour symmetry. Because this term does not commute with the generators of SU(6), the configuration mixing of different baryon states, predicted by Isgur and Karl [53], is the logical consequence. Having absorbed this term into the quark mass sector, which is known to be SU(3) broken anyway (but with apparently minimal consequences for the SU(3) baryon eigenstates), an important factor for potential SU(6) breaking has been eliminated. It is therefore important to notice the consistency when the F/D or z value does not indicate a substantial deviation from its SU(6) value.

Implementing SU(6) symmetry might seem inconsistent when we simultaneously observe polarised strange quarks inside a nucleon contributing to its magnetic moment (4.24). Is the NQM ruled out by this finding and by the suggested tiny quark spin sum (4.23)? After all, important particle properties, for example the $\Lambda - \Sigma$ mass difference, can be understood by applying QED and QCD corrections to the

NQM.

On the other hand there has never been a complete dynamical calculation of the NQM starting, for example from the QCD motivated effective Lagrangian which has constituent quarks, gluons and pions as basic degrees of freedom [86]. In this effective theory heavy particles are normally integrated out, and the first order approach involves only the light particles which are relevant at this particular scale between about 200MeV and 1GeV; that is between the quark confinement scale and the scale where the SSB of the chiral symmetry is expected to take place. In the next chapter we shall return to this model in more detail. Previously, the spin problem of the proton was seen as a high energy problem which disappeared when one reached the scale where the constituent quarks became the fundamental degrees of freedom. Now it is suggested that the spin problem is pushed further to the limit where the constituent quarks form a nucleon.

But because the quark spin structure analysis via the baryonic magnetic moments and via polarised deep inelastic scattering do not give any quantitative information about the absolute number of quarks and antiquarks of any flavour inside the proton one does not even have to worry about the most fundamental assumption underlying the NQM which is that the additive quantum numbers of three quarks which obey Fermi-Dirac statistics are sufficient for the classification of all known baryons.

The fundamental success of the NQM should be recovered by *any* more fundamental approach, but we can already summarise the following positive features from our analysis at the constituent quark level:

Especially encouraging is the new value of F/D or $z_{mag} \simeq z$. It should be the same at current and constituent quark level, and *is* actually found to be the same. Furthermore, it is found to be close to its SU(6) value.

The baryon wavefunctions remain eigenstates of the SU(3) flavour symmetry, and perhaps also of a somewhat modified SU(6) symmetry. It was not necessary for the analysis to break the spin-flavour-space symmetry assumption explicitly. The classical SU(6) breaking due to spin-spin hyperfine interactions is absorbed by the effective constituent quark masses.

Therefore the masses of the constituent quarks can vary with the energy of their environment, but the mass ratio defined by the SU(3) breaking parameter $\mathcal{S}$ in (4.16) can be determined very accurately (4.22). In fact, we regard the value of the mass ratio given by (4.22) as an inherent property of a constituent quark. Constituent quarks appear as physical and frame independent quantities, and their fields should be the fundamental degrees of freedom in an appropriate effective Lagrangian.

The future task will be the proper explanation of the baryons in terms of these constituent quarks.

5. Relativistic SU(6)?

5.1. Strangeness in the proton

Both the EMC experiment and the analysis of the baryon magnetic moments suggest a significant polarised strange quark contribution inside the proton. Several strangeness operators have been measured beforehand with different indications about the strangeness content of the proton. In the following, we review the still open question of the strangeness content inside the proton.

The smallest estimate of strangeness is suggested by deep inelastic neutrino scattering. The interpretation of the data, e.g. from the CDHS experiment [97], is given in the perturbative language of the parton picture and says that only $2.6 \pm 0.6\%$ of the nucleon's momentum is carried by s and $\bar{s}$ quarks:

$$\int_0^1 dxx(s(x) + \bar{s}(x)) = 0.026 \pm 0.006 \tag{5.1}$$

However, as we have mentioned already in chapter 3, these data do not constrain $\triangle s$. Different operators are involved in both measurements. Most general and simple inequalities among structure functions [98] that follow from unitarity can be written down in terms of their n^{th} moment M_n. Especially interesting are the following inequalities among the structure functions g_1 and F_1, which are valid because $F_1 > |g_1| > 0,\ 0 < x < 1$ and $|a+b| < |a| + |b|$:

$$|M_n(g_1)| = |\int_0^1 dx x^{n-1} g_1(x)| \leq \int_0^1 dx x^{n-1} |g_1(x)| \leq$$

$$\leq \int_0^1 dx x^{n-1} F_1(x) = M_n(F_1)$$

$$|M_n(F_1)| = \int_0^1 dx x^{n-1} F_1(x) \geq \int_0^1 dx x^{n-1} x F_1(x) = M_{n+1}(F_1) \tag{5.2}$$

Because CDHS does not measure $M_1(F_1)$, the first moment of F_1, but only the second moment $M_2(F_1)$, the matrix element $\bar{s}\gamma_\mu\gamma_5 s$ is not directly subject to any bound given by a different experiment.

Because of the generally non-trivial Q^2-dependence of the strange quark bilinear operators, it is in any context difficult to compare different data. Kaplan and Manohar [99] discussed the difficulty of extracting $\triangle\Sigma = \triangle u + \triangle d + \triangle s$ from data of elastic neutrino proton scattering. At very low energies the Z^o-boson measures the matrix element of the weak neutral current $\bar{u}\gamma_\mu\gamma_5 u - \bar{d}\gamma_\mu\gamma_5 d - \bar{s}\gamma_\mu\gamma_5 s$. The problem, however, lies in the extrapolation of the data of the elastic form-factors from a non-vanishing $q^2 = -Q^2$ to $\triangle\Sigma(q^2 = 0)$. A group at LAMPF is presently mounting an experiment to measure very low energy elastic neutrino scattering from liquid scintillators. Normally, the Q^2-dependence of the elastic form-factors is given in the dipole approximation, typified by the mass-scale $M_A \simeq 1.0GeV$ in case of the axial form-factor:

$$G_1^{(a)}(Q^2) = \frac{G_1^{(a)}(0)}{(1 + \frac{Q^2}{(M_A^a)^2})^2} \tag{5.3}$$

$M_A = 1.03 \pm 0.04 GeV$ has been measured in weak charged current processes which only involve the octet currents. Unfortunately, there are again no strict symmetry arguments for assuming that the singlet form-factor G_A^o of the weak neutral current has the same Q^2-dependence [99]. Only under the assumption that the singlet form-factors have the same Q^2-dependence do the existing BNL-data of the elastic νp scattering [11] support the EMC-result [4]:

$$\triangle u + \triangle d + \triangle s|_{\nu p} = 0.03 \pm 0.08 \tag{5.4}$$

Not least because of the meager statistics for the νp elastic scattering and the problem of the normalisation of the neutrino fluxes, this result is by no means as significant as the EMC-result. However, the latter might be taken in support of the assumption that the axial form-factors of elastic νp scattering are the same for the octets and singlet currents.

Another neutrino experiment, but in the deep inelastic region, has been proposed by Bass [100] who believes that the polarised analogy to the Gross-Llewellyn Smith sum-rule [101], viz.

$$\int_0^1 dx g_3^{\nu+\bar{\nu}}(x) = (\triangle q_{(v)})_o \tag{5.5}$$

measures anomaly-free the spin dependent valence quark distribution of

$$\triangle q_{(v)}(x) = (q^{\uparrow} - \bar{q}^{\uparrow})(x) - (q^{\downarrow} - \bar{q}^{\downarrow})(x) \tag{5.6}$$

This is not correct. As we have seen in the last chapter, $\triangle q_{(V)}$ corresponds to the magnetic moment operator σ_q, and an $s\bar{s}$ pair can give a non-vanishing contribution to the magnetic moment of the nucleon without smuggling an s-valence quark into the proton. Nevertheless, the measurement of the above sum-rule could provide the needed independent confirmation of the constituent quark spin sum as it is predicted through the magnetic moments (4.23/4).

An indication for the existence of non-naive (i.e. $s\bar{s}$) constituents inside the proton has been given for a long time by the πN sigma term analysis (see Decker et.al. [102] for a recent review, and Weise [103]). However, no general consensus about its value has been achieved so far.

The general σ-term is defined as:

$$\begin{aligned} \sigma_N^{ab} &= i \int d^3x < N|[A_0^a(\vec{x},0), \partial^0 A_0^b(0)]|N > \\ &= < N(p)|[Q^{5a},[Q^{5b},H(0)]]|N(p) > \end{aligned} \tag{5.7}$$

Hence the σ-term is the double commutator of the Hamiltonian density H with two axial charges Q^{5i}. The axial charges are part of the generators of the chiral symmetry which is assumed to be $SU(3)_L \times SU(3)_R$. The generators commute with the Hamiltonian as long as the chiral symmetry remains unbroken, yielding $\sigma_N^{ab} = 0$. We can thus replace H by the chiral symmetry breaking term. In the framework of QCD, the only possible chiral symmetry breaking term is the bare quark mass term:

$$\begin{aligned} H_B &= m_u \bar{u}u + m_d \bar{d}d + m_s \bar{s}s \\ &= c_0 u_0 + c_8 u_8 + c_3 u_3 \end{aligned} \tag{5.8}$$

with the standard definitions

$$
\begin{aligned}
c_0 &= \frac{1}{\sqrt{6}}(m_u + m_d + m_s) \qquad & u_0 &= \sqrt{\frac{2}{3}}(\bar{u}u + \bar{d}d + \bar{s}s) \\
c_8 &= \frac{1}{\sqrt{3}}(\frac{m_u + m_d}{2} - m_s) \qquad & u_8 &= \sqrt{\frac{1}{3}}(\bar{u}u + \bar{d}d - 2\bar{s}s) \\
c_3 &= \frac{1}{2}(m_u - m_d) \qquad & u_3 &= (\bar{u}u - \bar{d}d)
\end{aligned}
$$

H_B consequently transforms according to the $(3,\bar{3})+(\bar{3},3)$ representation of $SU(3)_L \times SU(3)_R$. Hence we obtain for the pion-nucleon σ-term, neglecting isospin violation, i.e. setting $m_u = m_d = m$

$$
\begin{aligned}
\sigma_N^\pi &= \frac{1}{2}(m_u + m_d) < N|\bar{u}u + \bar{d}d|N > \\
&= \frac{2m}{2m + m_s} < N|c_0 u_0|N > + \frac{m}{m - m_s} < N|c_8 u_8|N >
\end{aligned}
\tag{5.9}
$$

The ratio of the current quark masses $\frac{m}{m_s}$ is determined via PCAC in the meson sector [90]

$$
\frac{2m}{m_s - m} = \frac{m_\pi^2}{m_K^2 - m_\pi^2} \qquad \rightarrow \qquad \frac{m}{m_s} \simeq \frac{1}{25}
\tag{5.10}
$$

The nucleon mass shift due to the SU(3) breaking Hamiltonian $c_8 u_8$ is

$$
\triangle m_N =< N|c_8 u_8|N >
\tag{5.11}
$$

which is normally related to the general baryon octet mass splitting (Cheng& Li in [1]):

$$
\triangle m_N = \triangle m = m_\Lambda - m_\Xi \simeq -0.20 GeV
\tag{5.12}
$$

Under the assumption that the SU(3) singlet scalar operators and octet are not renormalised differently in the transition from current to constituent quarks, $< N|c_0 u_0|N >$ and $< N|c_8 u_8|N >$ can be related through

$$y = \frac{2 < N|\bar{s}s|N >}{< N|\bar{u}u + \bar{d}d|N >} \tag{5.13}$$

In the case that the renormalisation is different for singlet and octet, a strangeness component would automatically appear inside the nuleon. In the NQM, of course, any strangeness component is assumed to vanish, i.e. $y = 0$. In general, the σ-term can now be written

$$\sigma_N^\pi = \frac{3m_\pi^2(m_\Xi - m_\Lambda)}{2(m_K^2 - m_\pi^2)(1-y)} = \frac{\sigma_N^{\pi(y=0)}}{(1-y)} \simeq \frac{25MeV}{(1-y)} \tag{5.14}$$

In the above analysis SU(3) flavour symmetry in the baryon octet masses has been assumed to be broken by a power expansion of first order in $(m_s - m)$ around a common mass. Corrections of order $(m_s - m)^2$ are expected to be approximately $20\% - 30\%$ for SU(3) flavour symmetry breaking effects. Gasser and Leutwyler [104] calculated them in the framework of chiral perturbation theory with the result of a shift of $\sigma_N^{\pi(y=0)} = 25MeV$ to $\sigma_N^{\pi(y=0)} = 35 \pm 5MeV$.

However, our analysis of the baryon magnetic moments in an isospin invariant way, in which the SU(3) breaking pattern is systematically incorporated into the effective quark masses, suggests a different correction of $\sigma_N^{\pi(y=0)} = 25MeV$. The matrix element $< N|c_8 u_8|N >$ must be evaluated in the nucleon environment!

The SU(3) breaking effect shows up at the hadronic scale in the difference of the constituent quark masses m and m_s (i.e. in $S < 1$) but this difference has to be evaluated in a proton environment. Therefore we should use the 'rescaled' value for

$(m_s - m)|_N$ from equation (4.37), instead of the larger value $m_\Xi - m_\Lambda$. The result would be a dramatic reduction of $\sigma_N^{\pi(y=0)} = 25 MeV$ to

$$\sigma_N^{\pi y=0} = (13 \pm 5) MeV \tag{5.15}$$

This, on the other hand, fundamentally increases the gap between $\sigma_N^{\pi(y=0)}$ and its related 'experimental' value

$$\Sigma^{\pi N} = \sigma^{\text{'}exp\text{'}} = (64 \pm 8) MeV \tag{5.16}$$

which is however even less well determined. $\Sigma^{\pi N}$ is the product of the squared pion decay constant $f_\pi \simeq 92 MeV$ and the isospin even, on mass-shell, πN amplitude $\bar{D}^+$ with the Born term removed: $\bar{D}^+(t) = D^+(t) - \frac{g_{\pi N}^2}{M}$. As a consequence of the low-energy theorem of current algebra $\Sigma^{\pi N}$ and σ_N^π should be the same in the chiral limit of vanishing quark masses $m_q = 0$ [12]. The comparison of these two quantities, however, requires corrections due to the fact that σ_N^π is to be taken off mass-shell at $t = 0$. The corrections are expected to be of the same order as the sigma term itself, because its 'uncorrected' magnitude is already effected by the explicit breaking by the non-vanishing quark masses. Gasser, Leutwyler and Sainio [105,106] have analysed the dispersion relation of the form-factor $\sigma_N^\pi(t)$, with $t = (p - p')^2$ being the proton momentum transfer in πN scattering:

$$\sigma(t) = \sigma(0) + \frac{t}{\pi} \int_{4m_\pi^2}^{\infty} dt' \frac{Im\sigma(t')}{t'(t' - t - i\epsilon)} \tag{5.17}$$

where the spectrum $Im\sigma_N^\pi(t')$ is dominated by the low-mass s-wave $\pi\pi$ continuum. This 'soft' component of the meson cloud surrounding the nucleon gives rise to a very large mean squared scalar radius of the nucleon

$$< r_s^2 >= \frac{6}{\pi\sigma(0)} \int_{4m_\pi^2}^{\infty} \frac{dt}{t^2} Im\sigma(t) \tag{5.18}$$

The authors of [105] find that $< r_s^2 >^{1/2} \simeq 1.3 fm$ which has to be compared with the proton charge radius $< r_Q^2 >_p^{1/2} \simeq 0.86 fm$. The corresponding large slope of $\sigma(t) = \sigma(0)[1 + \frac{t}{6} < r_s^2 > + ...]$ implies the correction

$$\sigma(t = 2m_\pi^2) - \sigma(0) \simeq 15 MeV \tag{5.19}$$

and the corresponding value of $\Sigma^{\pi N}$ [106]

$$\sigma^{`exp'} = (45 \pm 5) MeV \tag{5.20}$$

This value, of course, is again closer to the value $\sigma_N^{\pi(y=0)} = 35 \pm 5 MeV$, which has been advocated in the framework of chiral perturbation theory [104]. Unless, however, the next corrections and the convergence criterion, are not calculated, there is no reason to conclude that $y = 0$ is the best solution. In contrast, the simple symmetry arguments which yield the environment dependence of the constituent quark masses suggest an even smaller value of $\sigma_N^{\pi(y=0)}$ giving rise to a quite large scalar strange form factor 'inside the nucleon' ($y \simeq 0.6$). Looking at the different charge and scalar radii of the proton the question is of course what 'inside the nucleon' means. All one can suspect at the moment is the presence of non-perturbative strangeness contributing to some properties of the nucleon. The reader who is interested in the evaluation of $< p|m_s\bar{s}s|p >$ should also consult the recent analysis by Jenkins and Manohar [107]. They calculate the leading non-analytic $m_s^{3/2}$ correction to the proton mass produced by meson loops and arrive so at a significant reduction of $< p|m_s\bar{s}s|p >$.

Further experimental evidence for a non vanishing strangeness content of the proton at large distances (low momenta) comes from the observation of OZI-forbidden processes [108,109], in which the vector ρ-meson and the tensor f'-meson (both almost pure $s\bar{s}$ states) couple directly to the proton. Predictions for further experimental tests of OZI-evading meson-baryon coupling are given by several groups [110].

In a recent publication Kaplan [111] claims to have found the strongest evidence so far for large strange matrix elements in non-strange particle states. His argument is based on the chiral quark model by Manohar and Georgi [86] and on the observation of the octet enhancement rule in $\triangle S = 1$, weak non-leptonic hyperon decays. Octet-enhancement stands for the fact that the terms in the Lagrangian

$$\mathcal{L}^{\triangle S=1} = \bar{u}\gamma^{\mu}(1-\gamma_5)s\bar{d}\gamma_{\mu}(1-\gamma_5)u + h.c. \tag{5.21}$$

which do not transform as (8,1) but as (27,1) under $SU(3)_L \times SU(3)_R$ are strongly suppressed by O(10). This rule is also known as the $\triangle I = \frac{1}{2}$ rule, because the operator which transforms as (27,1) gives rise to the suppressed $\triangle I = \frac{3}{2}$ amplitudes. The octet part can be parametrised in terms of the chiral baryon Lagrangian

$$\begin{aligned}\mathcal{L}_{\chi B} = \sqrt{2}G_F m_{\pi^+}^2 f_\pi(h_F tr\bar{B}[\xi^\dagger h\xi, B] \\ + h_D tr\bar{B}\{\xi^\dagger h\xi, B\} + \underline{27} + h.c.)\end{aligned} \tag{5.22}$$

with the definition $\xi = exp(i\pi^a T_a/f_\pi)$, where π^a is the pseudoscalar octet. The matrix $h = (T^7 + iT^6)$ changes an s to a d, and B is the baryon octet matrix

$$B = \begin{pmatrix} \frac{1}{\sqrt{2}}\Sigma^o + \frac{1}{\sqrt{6}}\Lambda & \Sigma^+ & p \\ \Sigma^- & -\frac{1}{\sqrt{2}}\Sigma^o + \frac{1}{\sqrt{6}}\Lambda & n \\ \Xi^- & \Xi^o & -\frac{2}{\sqrt{6}}\Lambda \end{pmatrix} \qquad (5.23)$$

When the operator (5.22) is fitted to s-wave non-leptonic hyperon decay data, the coupling constants h_F and h_D (corresponding to the F and D parameters) can be determined [112]

$$h_D = -0.58 \quad , \qquad h_F = 1.4 \qquad (5.24)$$

Kaplan points out the existence of the effective quark operator

$$O_s = 4(\bar{d}\gamma_\mu(1-\gamma_5)su\gamma_\mu(1-\gamma_5)\bar{s}) \qquad (5.25)$$

whose matrix elements between nucleons and pions do not vanish and are instead related by SU(3) symmetry to the measurable matrix elements from non-leptonic hyperon decays (5.21). In the NQM these matrix elements vanish. Kaplan argues that the large magnitude of the matrix elements $< n|O_s|p >$ and $< n\pi^o|O_s|p >$, given an s-wave $n\pi^o$ state, could only be cancelled by the (27,1) piece of O_s. But this piece is experimentally suppressed by octet enhancement.

However, Kaplan does not explain the exact matching conditions which have to be applied in order to parametrise the non-perturbative strong interaction effects. He does however, discuss in detail the renormalisation change from the weak scale $\mu \simeq M_W$ to $\mu \sim 1GeV$, the scale where the SSB of the chiral symmetry is expected to take place.

5.2. A dynamical description

This chiral symmetry breaking scale, usually denoted by $\Lambda_{\chi SB}$, has been specified by Manohar and Georgi [86] in the framework of their effective chiral quark model using 'naive dimensional analysis'. $\Lambda_{\chi SB} = 4\pi f_\pi$ is supposed to be the reference scale of the non-renormalisable effective Lagrangian which describes the dynamics below $\Lambda_{\chi SB}$. The goal of their work is the understanding of the success of the NQM (remember NQM stands for non-relativistic quark model). The effective Lagrangian involves quark, gluon and goldstone boson fields. The goldstone boson fields

$$B = \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^o + \frac{1}{\sqrt{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^o + \frac{1}{\sqrt{6}}\eta & K^o \\ K^- & \bar{K}^o & -\frac{2}{\sqrt{6}}\eta \end{pmatrix} \tag{5.26}$$

are defined in terms of a 3×3 matrix field $\Sigma(x)$ by

$$\Sigma = e^{2i\pi/f_\pi} \tag{5.27}$$

If in addition a ξ field is defined in terms of Σ

$$\xi = e^{i\pi/f_\pi} \quad , \quad \Sigma = \xi\xi \tag{5.28}$$

then the following combinations of pion fields transform as a vector and an axial vector

$$\begin{aligned} V_\mu &= \frac{1}{2}(\xi^\dagger \partial_\mu \xi + \xi \partial_\mu \xi^\dagger) \\ A_\mu &= \frac{1}{2}(\xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger) \end{aligned} \tag{5.29}$$

and ξ transforms under $SU(3)_L \times SU(3)_R$ as

$$\xi \quad \rightarrow \quad L\xi U^\dagger = U\xi R^\dagger \tag{5.30}$$

where U is defined by this transformation implicitly. The quarks, which are written as a set of colour triplet Dirac fermions

$$\psi = \begin{pmatrix} u \\ d \\ s \end{pmatrix} \tag{5.31}$$

transform as

$$\psi \rightarrow U\psi \tag{5.32}$$

The first few terms of the effective Lagrangian which would be invariant under chiral SU(3) without the mass term are

$$\mathcal{L} = \bar{\psi}(iD_\mu + V_\mu)\gamma^\mu\psi + g_A\bar{\psi}A_\mu\gamma^\mu\gamma_5\psi - m\bar{\psi}\psi$$
$$+\frac{1}{4}f_\pi^2 tr\partial^\mu\Sigma^\dagger\partial_\mu\Sigma - \frac{1}{2}trF_{\mu\nu}F^{\mu\nu} \tag{5.33}$$

where $D_\mu = \partial_\mu + igG_\mu$ is the usual covariant derivative, and $F_{\mu\nu} = \partial_\mu G_\nu - \partial_\nu G_\mu + ig[G_\mu, G_\nu]$ is the gluon field strength tensor.

The constituent quark mass m is produced by chiral symmetry breaking and is proportional to the vacuum expectation value of the chiral condensate:

$$m \simeq g^2 \frac{< \bar{\psi}\psi >}{q^2} \tag{5.34}$$

m becomes 'soft', i.e. it falls off with Q^2 only at momenta large compared to $\Lambda_{\chi SB}$. Embedding this expression for m into a hadronic environment, small fluctuations of m due to a different effective coupling g and due to a slightly changed value of $< \bar{\psi}\psi >$, seem to be quite likely and do not alter the dynamics of the chiral quark model to first order. Including hyperfine splitting effects already into the definition of the effective coupling can change the dynamical constituent quark mass in the right direction.

In terms of the Nambu Jona-Lasinio (NJL) model [95] a more detailed study of the dynamics of the SSB of chiral symmetry and the corresponding dynamical constituent quark mass generation can be given. (See e.g. Vogl and Weise for a pedagogical introduction to all the potential and the subtleties of the NJL model [113]) Bijnens, Bruno and de Rafael [114] found in their study of the NJL model that the chiral quark model is a special case of the NJL model, but the connection is non-trivial. Probably the main distinction between the chiral quark and the NJL model is the fact that in the chiral quark model all but pion interactions with the quark fields are neglected. Thus different phenomenological predictions (e.g. pion-meson mixing) are given in those models.

On the other hand, we find that in the long wavelength and low-energy limit of the NJL model all gluon degrees of freedom are frozen into effective pointlike interactions between quarks. This is a reasonable assumption for the low-energy approximation of QCD, where the global symmetries of QCD are thought to be more important than confinement effects. However, essential gluon features like the QCD $U_A(1)$ anomaly have to be matched with effective quark interactions only.

It is still under debate [115], if pointlike quark interactions do fulfill this task. 't Hooft [116] introduced a determinant interaction in order to break the $U_A(1)$ symmetry of the $\mathcal{L}_{QCD}$ via instantons. With three flavours, the effective six-point quark interaction

$$\mathcal{L}_6 = \frac{1}{2} G_D(det[\bar{q}_i(1+\gamma_5)q_j] + det[\bar{q}_i(1-\gamma_5)q_j]) \tag{5.35}$$

uniquely breaks $U_A(1)$ without breaking $SU(3)_L \times SU(3)_R \times U_V(1)$. This symmetry is also respected by the four-point interaction

$$\mathcal{L}_4 = G_s \sum_{a=0}^{8} [(\bar{q}\frac{\lambda^a}{2}q)^2 + (\bar{q}i\gamma_5\frac{\lambda^a}{2}q)^2] + \dots \tag{5.36}$$

For simplicity, only the leading scalar-pseudoscalar term is written down. The total NJL Lagrangian still involves the kinetic and the current mass terms of the quark fields:

$$\mathcal{L}_{NJL} = \bar{q}(i\partial^\mu\gamma_\mu - m)q + \mathcal{L}_4 + \mathcal{L}_6 \tag{5.37}$$

However, one has to keep in mind that many studies, e.g. [114], are still done in the safe $N_c \to \infty$ limit where the effects of the gluon anomaly disappear and $G_D = 0$. On the other hand, with $G_D \neq 0$ one observes interesting effects of flavour mixing, which have been studied by several groups.

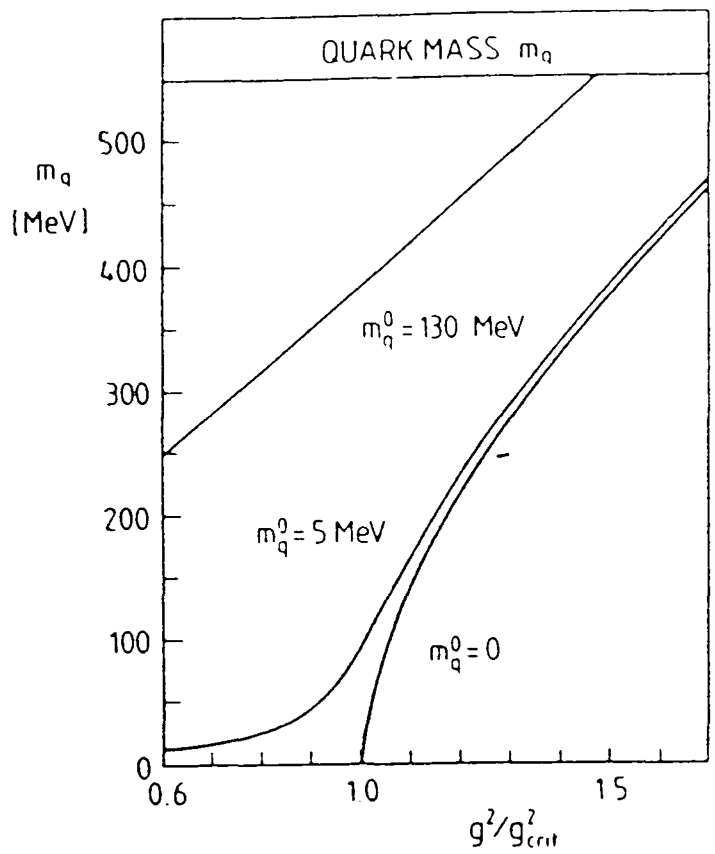
As consequence of the spontaneous chiral symmetry breaking due to a strong effective coupling (after exceeding a critical value) even massless current quarks acquire a dynamical effective constituent quark mass in the Hartree (mean field) approximation [113]. By 'closing the quark-loop' the multiple fermion interaction can be linearised, and the following mass gap equations are obtained (see ref.[113] for a detailled description):

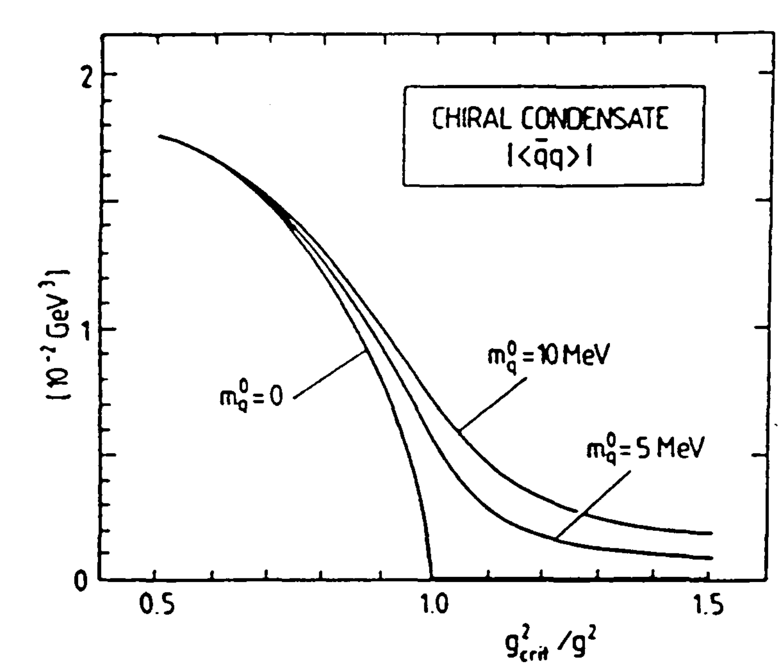
$$
\begin{aligned}
m_u &= m_u^o - G_s < \bar{u}u > -G_D < \bar{d}d >< \bar{s}s > \\
m_d &= m_d^o - G_s < \bar{d}d > -G_D < \bar{u}u >< \bar{s}s > \\
m_s &= m_s^o - G_s < \bar{s}s > -G_D < \bar{u}u >< \bar{d}d >
\end{aligned}
\tag{5.38}
$$

where m_q^o is the current quark mass. For each flavour the quark condensate is given by

$$
< \bar{q}q >= -4iN_c \int^{\Lambda} \frac{d^4p}{(2\pi)^4} \frac{m_q}{p^2 - m_q^2 + i\epsilon} \tag{5.39}
$$

This integral is usually regularised by introducing a characteristic cutoff $\Lambda \sim 1GeV$. Ultimately, all free parameters which apppear in $\mathcal{L}_{NJL}$ have to be fitted to observable quantities. Interesting with respect to the environment dependent constituent quark mass is the behaviour of the dynamical quark mass m_q and of the chiral condensate $< \bar{q}q >$ as a function of g^2 and g^{-2} respectively. Here $g^2 = (G_S + G_D < \bar{q}q >)\Lambda^2$ is the squared dimensionless coupling constant [117]:





Already a small shift in the effective coupling (e.g. due to hyperfine interactions) could change the dynamical quark mass of all flavours by a few percent.
It is not yet clear whether the observed polarised strangeness component of the proton and its quark spin structure can be described successfully in the framework of the NJL model, and whether the flavour mixing six-point interaction is necessary. Some features have been calculated starting from $\mathcal{L}_{NJL}$ and appear quite attractive [117][118][119] [120]

The matching conditions for the free parameters in an effective field theory are basically all put in by hand. In the chiral quark model the effective strong coupling constant

$$\alpha_s(Q^2 \to \Lambda_{QCD}^2) \sim 0.3 \qquad (5.41)$$

is not greater than one. This is in contrast to what one might have expected from the definition of α_s in (1.8). The reason why α_s can have a small effective value (5.41) is given by the different nature of the quark fields involved in the effective Lagrangian. This phenomenon has been explained by Manohar and Georgi also in the following intuitive way [86]:

The chiral symmetry breaking occurs at a scale $\Lambda_{\chi SB} \sim 1GeV$ because $\alpha_s(1GeV^2)$ becomes so large that $\bar{\psi}\psi$ gets a non vanishing vacuum expectation value. Thus the large value of α_s destabilises the original vacuum with respect to the one where the chiral symmetry is spontaneously broken. Since the reason for the instability was a large α_s, it is clear that this broken-symmetry vacuum will have a smaller value of α_s.

The value of α_s (5.41) in the effective theory has been matched to the baryon spectrum by looking at mass shifts which are caused by QCD corrections. The small value of α_s suggests that the baryons are quite loosely bound objects. Manohar and Georgi argue that the value of α_s (5.41) is so small that the effective quarks are indeed non-relativistic objects. This argument would be weakened by Weinberg's analysis of the low-energy theorems in the framework of the chiral quark model. He determines $g_A = 1$ and $\kappa_q = 0$ for the constituent quarks. In this case, however, the constituent quark mass inside the proton (4.38) becomes so small

$$m_u|_N \simeq 250MeV \qquad (5.42)$$

that the nonrelativistic approximation would require a smaller value of α_s.

Assuming the NQM as the valid underlying theory at the hadronic scale, the matching condition of the weak coupling g_A in the effective Lagrangian requires the reduced effective value [86]

$$g_A = 0.75 \tag{5.43}$$

due to the difference between the NQM value of $\frac{g_A}{g_V}$ and its experimental value, as measured in the neutron β-decay. In this case, the constituent quark mass is about $m_u|_N \simeq 330MeV$.

The value of g_A is automatically reduced in the NJL model, where the pions are allowed to mix with the a_1 axialvector meson [113][114]. Apparently this effect alters the value of $g_A = 1$ (as calculated by Weinberg) to $g_A < 1$. A clarification for this difference is outstanding, and should be related to the problem of how the pion states in the chiral quark Lagrangian are to be identified with the observed pion states. In the chiral quark model the mixing between the fundamental π^a and the additional $q\bar{q}$ bound state with the same quantum numbers yield the separation between the massless pion and the next pseudoscalar state. Its mass is very heavy and therefore outside the effective theory. Thus, double counting can be avoided. The NJL model gives more justice to the 'external' degrees of freedom by including the vector and axialvector mesons etc. into the analysis. In the NJL-model all mesons are taken as quark-antiquark composites, even the light pseudoscalars.

In the chiral quark model and in the NJL model, the anomalous magnetic moment κ_q of the constituent quark has been found to be negligible [18][113]. This is an essential observation. The spontaneous chiral symmetry breaking generates a significant dynamical mass for the constituent quark. Nevertheless it remains a pointlike object (i.e. a bare Dirac particle) below $\Lambda_{\chi SB}$. The small α_s (5.41) does not suggest that a constituent quark is bound stronger than a current quark. One

can suspect that a constituent quark enjoys a similar asymptotic freedom in its s-wave state as a current quark experiences only at much higher energies (1.8).
But why don't we see any clear peaks or resonances due to these pointlike (and only 'weakly' strongly interacting) constituent quarks in the lepton nucleon scattering experiments? Similar peaks can be observed, for example, in a nucleus when the incoming e^- interacts elastically with one quasifree nucleon.
The answer to this may be that the smearing out of the quark distribution function at a Bjorken-x value of $x_{Bj} = 1/3$ is due the fact that the absolute number of constituent quarks and antiquarks is not necessarily fixed to three. So, the flavour of the hit constituent quark can either be up or down or strange.

Nevertheless, the observation of pointlike $s\bar{s}$ constituent quarks inside a nucleon via its magnetic moment is still somewhat disturbing, because we do not yet know the flavour-spin wavefunction analogy to the NQM one. Of course, the Pauli principle and the classification scheme of the baryon states should be obeyed. Also the remaining success of the NQM, i.e. mainly the explanation of all the higher resonances [121], should be incorporated into any new model.

SU(6) is definitely a reasonable underlying symmetry for the hadronic classification. Because the strong interaction does respect the SU(3) flavour symmetry and does not display any spin dependence to first order either, the greater SU(6) symmetry should hold at least at the approximate level used in the successful NQM.

Up to second order SU(3) breaking effects (e.g. due to a 'lower' and mass dependent quark spinor component), most quark models assume the F/D ratio (or z) to be the same at current and constituent quark level. Therefore, the more accurately determined constituent quark value $z_{mag} = 0.235 \pm 0.002$ is used in the subsequent final analysis of the EMC experiment. Because of the uncertainty about the magnitude of g_2 in the β-decays, the current quark value $z = 0.23 \pm 0.04$ is less well determined. (Karl specified for the sake of $z_{mag} = z$ the class of shell models with one-body quark and antiquark operators, where this identity must hold. So, at least in the framework of these models the following conclusions are compelling) From the experimental data of $\int_0^1 dx g_1^p(x) = 0.126 \pm 0.010 \pm 0.015$ we find the following spin components in parton language:

$$\begin{aligned} \triangle u - \frac{\alpha_s}{2\pi}\triangle G &= +0.77 \pm 0.05 \\ \triangle d - \frac{\alpha_s}{2\pi}\triangle G &= -0.49 \pm 0.05 \\ \triangle s - \frac{\alpha_s}{2\pi}\triangle G &= -0.31 \pm 0.05 \end{aligned} \tag{5.44}$$

Now we take the 'ideal matching' $g_A = 1$ for a constituent quark (assuming that no corrections make g_A significantly smaller than one anyway). Then the analysis of the baryon magnetic moments (4.24) suggests the following constituent quark spin components inside the proton:

$$\begin{aligned} \sigma_u &= +0.81 \pm 0.01 \\ \sigma_d &= -0.45 \pm 0.01 \\ \sigma_s &= -0.26 \pm 0.01 \end{aligned} \tag{5.45}$$

In the case of the magnetic moments no gluons contribute with their polarisation. The comparison of the two equations (5.44) and (5.45) is possible, if we strictly assume the spin structure to be scale independent after it has been written in this form. If we combine these equations we find that

$$\triangle \bar{q} = -0.02 \pm 0.03 + \frac{\alpha_s}{4\pi}\triangle G \tag{5.46}$$

where the antiquark polarisation is defined by $\triangle\bar{u} = \triangle\bar{d} = \triangle\bar{s} = \triangle\bar{q} = \bar{q}^{\uparrow} - \bar{q}^{\downarrow}$. Unfortunately, it is not straightforward to argue with respect to the evolution equations of perturbative QCD, that the antiquark polarisation contributes necessarily with the same sign to the proton spin as the gluon polarisation. Here we compare constituent and current quark spin.

Anyway, the sign identity of $\triangle\bar{q}$ and $\triangle G$ is not sufficient for explaining why $\triangle s < 0$ and $\sigma_s < 0$. Of course, one has to keep in mind that any antiquarks contributing to the baryon magnetic moments have effective dynamical constituent masses.

On the other hand we can use Weinberg's theoretical explanation and our 'experimental' observation to argue for constituent quarks being pointlike particles with no intricate spin structure. We set $\triangle\bar{q} = \triangle G \simeq 0$ and approximate the values of the equations above to:

$$u^{\uparrow} - u^{\downarrow} \equiv +\frac{3}{4} \quad , \qquad d^{\uparrow} - d^{\downarrow} \equiv -\frac{1}{2} \quad , \qquad s^{\uparrow} - s^{\downarrow} \equiv -\frac{1}{4} \tag{5.47}$$

What is the explanation for such a result? Does this result make the situation worse? What is the consequence of (5.47)? Let us consider this carefully.

A first affirmative observation is the fact that all values in equation (5.47) are shifted equally from the naively expected values:

$$u^{\uparrow} - u^{\downarrow} = +1 \quad , \qquad d^{\uparrow} - d^{\downarrow} = -\frac{1}{4} \quad , \qquad s^{\uparrow} - s^{\downarrow} = 0 \tag{5.48}$$

The NQM picture could be recovered, if there was a non-perturbative mechanism, like a topological gluon anomaly in the flavour singlet channel [122]. This had to change exclusively the values of the quark spin components by $-\frac{1}{4}$. The idea of instantons causing non-trivial admixtures to bilinear quark operators has not yet been proven to be calculable [122][123][124].

The reason why the NQM picture cannot be recovered trivially is the SU(3) invariant reduction of the total quark spin sum. However, the SU(6) symmetry is not directly jeopardised in equation (5.47). This may be verified by calculating F/D (or z) from these numbers, which are indeed the SU(6) characteristic numbers. It is important to note that all the experiments which we looked at do not alter significantly the SU(6) symmetry as it is characterised by its $F/D = 2/3$ (or $z = 0.2$) value.

Without reference to gluons which are charge conjugation invariant, it is possible to provide an intuitive picture for the observed phenomena:

Starting from an effective low-energy theory (e.g. the NJL-model or the chiral quark model) one is able to investigate the impact of the meson cloud surrounding a constituent quark. In the chiral quark model it is necessary, for this purpose, to take into account the additional assumption of loosely bound quarks. Quarks have to interact directly with the fundamental pseudoscalars instead indirectly through nucleon-pion interaction. In the NJL-model, on the other hand, the constituent quarks are not confined and interact with all meson states (i.e. $q\bar{q}$ composites).

A qualitative understanding of the appearance of negative polarised s quarks with respect to positive polarised u quarks and hardly polarised $\bar{s}$ quarks is possible in both models by looking at the effect of the kaon cloud surrounding the u quark. In the virtual process

$$u \rightarrow sK^+ \rightarrow u \tag{5.49}$$

the pseudoscalar K^+ is emitted in a p-wave. Angular momentum conservation demands the s quark to be polarised antiparallel to the initial u spin. Because the $\bar{s}$ is in a spinless K^+, it is not polarised through this process. Quantitative predictions [125], which go beyond the $N_c \to \infty$ approximation or beyond the $1+1$ dimensions in the context of the Gross-Neveu model [126], are still outstanding.

Naturally, one may expect that other experiments do not provide the same evidence for the quark spin components in eq. (5.47). However, on the other side, most spin related high energy physics experiments underline indeed an insufficiency of the NQM in this particular respect.

For example, the data of the inclusive polarized baryon production cannot be explained in the NQM, as it was pointed out e.g. in ref.[9].

Scattering a proton off a heavy nucleus target (e.g. Beryllium) may produce polarized Λ, Σ or Ξ particles or antiparticles, and unobserved further particles.

In one of the latest attempts to understand the inclusively polarized Λ production, the observed spectrum could be explained in terms of the binary reaction $\pi p \to K\Lambda$, but not in terms of a more basic quark pion process [127]. The fact that the inclusively produced $\bar{\Lambda}$ is not polarized, suggests that the Λ is polarized in this binary reaction due to the negative polarized strangeness in the proton. In contrast, the anti-strangeness in the proton and (therefore?) the produced $\bar{\Lambda}$ particle have no polarization. The fact that $\bar{\Lambda}$ is produced without polarisation does not support the idea of the creation of separate constituent $s\bar{s}$ pairs, as proposed in the Lund model [128]. In this model, an $s\bar{s}$ state out of the vacuum has to be formed and is measured to have an angular momentum (via a nonvanishing p_T of the produced Λ). Therefore it is expected that the $s\bar{s}$ pair has an internal spin opposite to $L_{s\bar{s}}$, that gives rise to the polarisation of the Λ. This implies that an inclusively produced $\bar{\Lambda}$ has to be polarised as well.

In general processes which involve hadron-hadron collisions are complex and a plausible phenomenological explanation cannot guarantee the validity of the expected underlying mechanism. A complete dynamical mechanism should be given instead.

Finally the question arises why the quark spin sum inside the baryons is compatible with zero, and what are the consequences for the solution of the $U_A(1)$ problem. An argument by Jaffe [129] proposes that a vanishing flavour singlet axial vector matrix element $a_0 = 0$ could trivially solve the $U_A(1)$ problem.
We found in the second chapter that it is impossible to calculate a_0 via the η' coupling constants, because of the gluon anomaly. However, with a vanishing constituent quark spin sum being compatible with the data (independent of the gluon anomaly), the conserved chiral charge Q_5^0 can be identified over the constituent quarks with the vanishing quark spin sum in equations (5.47). Originally, Q_5^0 is defined over the non-observable (gauge variant), but conserved, current $A_\mu^0 - n_f k_\mu$. If $Q_5^0 = 0$, there are naturally no parity doublets in the hadronic spectrum. The $U_A(1)$ symmetry is not spontaneously broken and the η' cannot be seen as a mixture with any 'would be' NG boson of the flavour singlet axial vector channel.

A vanishing quark spin sum may have the following explanation:
A general problem of a frame-independent relativistic extension of the NQM is the fact that the NQM has a spin-1/2 and a spin-3/2 representation and is therefore rest frame dependent [130]. However, in a truly relativistic quark model the internal symmetries of spin-flavour have to commute with the space-time symmetry.
On the other hand, the obvious spin and flavour independence of QCD makes a greater SU(6) symmetry plausible. This symmetry, characterised by its typical F/D (or z) value, is not seriously questioned by experiment. However, its basic representation is questioned by the observation of a vanishing quark spin sum.
If also the decuplet states are discovered to have a vanishing internal quark spin sum, the flavour decuplet and octet representations might be part of a single 56-dimensional irreducible representation of SU(6). The internal quark spin symmetry

does not mix anymore with the external space symmetry, because the total quark spin sum is the same in flavour decuplet and octet states. A derivation of SU(6) from QCD becomes feasible.

Unfortunately, it is impossible to construct the decuplet states without knowing the spin-flavour symmetric wavefunction of the proton. Therefore, it is not clear how the measured magnetic moment of the Ω should be understood.

The isoscalar Ω^- is about three times the magnetic moment of the isoscalar Λ:

$$\mu_{\Omega^-} = -1.94 \pm 0.17 \pm 0.14 \tag{5.50}$$

The crucial question is, whether SU(6) symmetry is consistent with the experimentally verified relation

$$\mu_{\Omega^-} \simeq 3\mu_{\Lambda^\circ} \tag{5.51}$$

Since the quark spin sum inside the Λ vanishes

$$s^{\uparrow}(\Lambda) - s^{\downarrow}(\Lambda) = +\frac{1}{2} \quad , \qquad u^{\uparrow}(\Lambda) - u^{\downarrow}(\Lambda) = d^{\uparrow}(\Lambda) - d^{\downarrow}(\Lambda) = -\frac{1}{4} \tag{5.52}$$

the magnetic moment of the Ω does not necessarily suggest that it is based on the magnetic moments of three strange constituent quarks, which carry all the spin in the NQM. Nevertheless, the latter explanation is not excluded.

Naturally the Ω^- is not identical with a particle which is three times the Λ^0, and (5.51) might not be a relation of the SU(6) symmetry. Taking into consideration the heavy masses of the decuplet resonances, it is questionable if their explanation relies on a relativistic quark model.

However, after we have argued for processes in which the kaon cloud interacts directly with the constituent quarks, the decuplet resonances should be included into the effective theory. The mass difference between the octet and decuplet states is about the mass of a pseudoscalar meson. In the framework of effective quark theories the incorporation of decuplet states into the calculation of dynamical processes is therefore favoured [131]. The kaon cloud gives rise to negative polarised strangeness in the nucleon (5.49) and can also be responsible for negative polarised up and down quarks inside a polarised Ω^-. The ultimate confirmation about a vanishing quark spin sum inside the decuplet states can be given by measuring the spin dependent structure function $\int g_1^{\Omega^-}$.

In the context of the heavy particle effective theory, the baryon is taken as a heavy particle interacting with the light pseudoscalar Σ-field (5.27). Carone and Georgi [132] recently analysed non-leptonic decays of the octet and decuplet baryons in the framework of the chiral Lagrangian for $J^P = \frac{1}{2}^+$ octet and $J^P = \frac{3}{2}^+$ decuplet baryons. They found by looking at the non-leptonic Ω^- decays no simple SU(6) transformation law for the meson-baryon interaction. According to the authors the reason for this failure lies in a probable misunderstanding of the $\triangle I = \frac{1}{2}$ rule. This rule is obviously not needed for the description of the Ω^- decay.

A more general question might be addressed by trying to understand the non-leptonic Ω^- decay at quark level. In the usual relativistic extension of the NQM, where the particles are all boosted along the z-axis, the resulting $SU(6)_W$ symmetry transforms antiquarks differently from quarks. In this reference system these well-known non-leptonic decays are not forbidden.

However, the problem with this explanation is the possibility of π being a fundamental particle, i.e. a NG boson which is massless in the chiral limit $m_\pi^2 = 0$. The $SU(6)_W$ symmetry is however a simple extension of the NQM, where the π is seen as a $q\bar{q}$ boundstate. A relativistic invariant SU(6) theory has to explain these decays at the quark-pion level. This is possible if the quark spin sums in the flavour decuplet and octet states are equal. With this final speculation I would like to argue, that the observation of a nearly vanishing quark spin sum at constituent quark level may lead us to interesting questions for future research.

6. Conclusion

Several experiments have provoked the re-examination of the Nonrelativistic Quark Model (NQM). Most of them indicate directly that the weak part of the NQM is related to the incorrect description of the spin structure of the baryons.

In this thesis we have studied first the parton model interpretation of the polarised deep inelastic scattering experiment of the European Muon Collaboration (EMC). Equation (2.25) displays the almost vanishing combination of the quark spin sum $\triangle\Sigma$ and the gluon polarisation factor $-n_f\triangle\Gamma$. $\triangle\Gamma$ is proportional to $\alpha_s(Q^2)\triangle G(Q^2)$ which scales in leading order. It has been emphasized that $\triangle G$ is well defined (without ambiguity) through observable two-jet processes which belong to the graph in (2.13). The parton model is physically most suitable for comparison of current and constituent quarks because the transverse momentum k_T of the partons are low.
Because we do not observe a dense degenerate spectrum of baryons with different angular momentum, a large $\triangle G$ is not plausible for the groundstate. The consequence of this is a negative polarised strangeness in the proton, i.e. $\triangle s < 0$. This leads us to the *real* problem formulated by Altarelli: What is the reason for such a large discrepancy between the current and constituent quarks with respect to their spins?

The quantum field theoretical calculation of a_0, the axial vector flavour singlet matrix element, has recently been proven by Shore and Veneziano to be beyond current techniques, due to the gluon anomaly (1.13) and its nonperturbative effects, e.g. (2.37). Attempts to relate the η' decay to a_0 are wrong due to the Renormalisation Group (RG) variant parameter '$f_{\eta'}$' which has to be replaced by the unmeasurable constant F. In every conceivable relation where F appears (e.g. in the generalised $U_A(1)$ Goldberger-Treiman relation) a glue vertex with highly nontrivial RG properties enters as well (e.g. in form of the g_{GNN} glue-nucleon coupling constant). From this perspective it can be seen as success to have chosen a different approach to the spin problem.

The SU(3) flavour symmetry plays a crucial role when the spin structure of the baryons is analysed. The pattern of SU(3) breaking and its consequences for the Quark Model have to be understood if one is interested in the Cabibbo theory of the weak currents or in the question if baryons are still SU(3) eigenstates despite its breaking.

Also in order to test the unitarity of the Kobayashi-Maskawa matrix (3.3), the flavour octet axial vector currents are parametrised by the two characteristic SU(3) couplings F and D (3.6). Those are obtained from the $\frac{g_A}{g_V}$ form factor ratios of different baryon β decays. However, a fit to all β-decays does not yield an exact value (3.8). This indicates the presence of SU(3) breaking effects.

Because g_A is not protected via the Ademollo Gatto theorem against corrections stemming from baryon mass differences, we expect larger corrections naturally for the semileptonic hyperon decays. These corrections are not well matched by relativistic corrections (3.11) which give larger modifications to the form factors of the nucleon. A different scheme for corrections arising from SU(3) breaking involves baryon configuration mixing. In this case, the theoretical predictions by Zenczykowski (3.12) are close to the experimental data, but the author himself comments that 'arbitrary parameters are tailored to our needs'.

The first measurement of the induced second class form factor g_2/f_1 (3.17), however, produces a shift in g_1/f_1 proportional to the energy transfer in the β decay. This is in agreement with the fact that strong interaction effects can modify the hyperon decay form factors, because the V-spin is not as good as the Iso-spin symmetry. This observation does not imply the existence of second class currents. The only weak hadronic currents are the standard left-handed quark currents which are classified as first class currents. Already in a NQM approximation (3.20) the ratio $g_2/g_1 \simeq 0.73$ is not negligible.

Taking the measurement of g_2/f_1 into account, a new value for F/D can be calculated. This value $F/D = 0.73 \pm 0.09$ (3.27) can be more transparently expressed (3.23) by the ratio $z = \frac{a_3}{\sqrt{3}a_8} = 0.23 \pm 0.04$ (3.26). $z = 0.2$ is the SU(6) characteristic value. This new value of F/D has a mean value significantly above the old value $F/D = 0.58 \pm 0.05$ (3.8). Future independent measurements of g_2/f_1 and consequently g_1/f_1 in the $\Lambda \rightarrow p$ (3.29) and the $\Xi^- \rightarrow \Lambda$ (3.28) decays have to confirm these new values, lending further support to the suggested pattern of SU(3) breaking.

With such a new value of z no 'decent burial of SU(3)' (Lipkin) is necessary and the wavefunctions of the baryons remain SU(3) eigenstates. The SU(3) breaking is fully taken into account by unequal quark masses and the corresponding shift in g_1/f_1.

Assuming the validity of SU(3), the quark spin contributions to the proton spin are directly evaluated from the EMC result over the F/D (or z) value.

In this sense, the size of $\triangle s$, for example, depends crucially on the exact value of F/D (or z) (3.34). However, the observation of an extremely smal value for a_0, and consequently the existence of the spin problem, is not especially sensitive to F/D (or z) (3.36). Also the validity of QCD does not depend on any specific value of F/D. In this respect Close's defence of QCD against Preparata's and Soffer's critique is misleading. The bound on $|\triangle s|$ in (3.33) is not given by the neutrino scattering data (5.1) as it can be derived from the structure function inequalities in (5.2).

Because the anomalous dimension of the flavour octet axial vector currents vanishes ($dim[\partial^\mu A^j_\mu] < 4$), no large scale dependence of their matrix elements a^j is to be expected. This is the reason why the Bjorken sum-rule (2.22) holds and why the constituent quarks should match the current quarks with a common F/D.
It is a very important proof of consistency that the analysis of the baryon magnetic moments yield very accurate values for $F/D = 0.743 \pm 0.005$ (4.21) and $z_{mag} = 0.235 \pm 0.002$ (4.20) in perfect agreement with (3.27) and (3.26).

An independent analysis of the magnetic moments became necessary after the EMC result could not be understood at the current quark level. In addition to this dilemma, new magnetic moment data showed an up to 37.76% difference to their NQM prediction. According to Weinberg, a constituent quark is a bare Dirac particle with no anomalous magnetic moment $\kappa_q = 0$ and $g_A = 1$. However in this case, the discrepancy between a theoretical model prediction and the data of the baryon magnetic moments has to be much smaller.

Initially the parameter free prediction of $\frac{\mu_p}{\mu_n} = -\frac{3}{2}$ (1.2) has been seen as a big success of the NQM. However, this prediction was criticised by Bell for its frame dependency. Bell argued that the magnetic moment measurements involve a small but non vanishing momentum tranfer and should therefore be understood in a relativistic extension of the NQM.
Another question related to $\kappa_q = 0$ is the fact that in a harmonic oscillator potential κ_q vanishes only if the confinement radius $R \to \infty$ (1.5). However, because the

picture of confined quarks is not yet fully understood, we might adopt the idea that also constituent quarks are asymptotically free particles. Being only in second order subject to strong interactions, constituent quarks, gluons and pions can be seen as fundamental degrees of freedom in the effective Lagrangian $\mathcal{L}_{\chi QM}$ (5.33) of the chiral constituent quark model, with an effectively small coupling α_s (5.41).

However, even without invoking special constituent quark models, it is legitimate to start with the analysis of the baryon magnetic moments with the general ansatz (4.7).

Under the assumption of baryons being in SU(3) eigenstates, the set of equations (4.10) should recover the NQM or any bag model, if they are indeed the correct models for the description of the magnetic moments.

Due to the symmetric structure of (4.10), the free parameters are not sufficient for the solution of (4.10), as it is evident by the failure of the sum rules (4.11) and (4.12). However, a numerical analysis of (4.10) is feasible and Karl's minimalisation procedure of the error (4.14) yield a stable solution. The quark spin sum (4.15) is compatible with the EMC result, and the value of F/D (or z), which Karl finds, is in perfect agreement with the β-decays (3.26). Karl, who did not know this new value (3.26), could not explain the origin for the inconsistency which one had with the old value (3.8).

The systematic evaluation of (4.10) is the major advantage of our analysis of the baryon magnetic moments. Because isospin symmetry is much better than SU(3), (4.10) has to be analysed in an isospin invariant manner (4.17). Three pairs of equations from (4.10) are sufficient for the calculation of $F/D = 0.743 \pm 0.005$ (4.21) or ($z = 0.235 \pm 0.002$ (4.20)), for the calculation of the SU(3) breaking parameter $\mathcal{S} = \frac{m_u}{m_s} = 0.653 \pm 0.024$ (4.22), and for the calculation of the constituent quark spin sum $\sigma_u + \sigma_d + \sigma_s = (0.086 \pm 0.026)(\sigma_u - \sigma_d)$ (4.23).

This most accurate determination of $\mathcal{S} = \frac{m_u}{m_s}$ from the baryon flavour octet data is a clear signal for constituent quarks being observed, independent of the nonrelativistic frame. However, constituent quarks are depending on their environment. It follows from our isospin invariant analysis that all constituent quark masses scale (4.25),(4.30) and (4.32) with the mass of the baryon they are embedded in.
Usually, some baryon mass differences are explained through a flavour dependent two body interaction (4.34). Now it looks as if the effective constituent quark mass includes this interaction (4.36). This is a gain of simplicity. The SU(6) breaking hyperfine interaction does not result in a random mixing of baryon states, but is systematically absorbed to first order by environment dependent fluctuations of the effective constituent quark mass.

The concept that constituent quarks are fundamental physical quanta is not yet generally accepted. Our analysis of the magnetic moments suggests that constituent quarks are, indeed, observable as bare Dirac particles in experiments which measure quark properties below the chiral symmetry breaking scale. Above this scale of about 1GeV, quarks are perceived as current quarks with the characteristic mass ratio $m_u : m_d : m_s = 0.5 : 1 : 20.1$.
This mass ratio is significantly different from the measured constituent quark mass ratio $m_u : m_d : m_s = 1 : 1 : \mathcal{S}^{-1}$. Having only less accurate information about the absolute quark masses, the appearance of $s\bar{s}$ pairs inside a nucleon must , in principle, not be more of a surprise at the constituent quark level than it is at the current quark level.
It is important to note that the observation of strange and anti-strange constituent quarks inside a nucleon does not cast any doubt on the valence quark structure of the baryons. For the determination of the valence quark distribution inside a baryon the antiquark distribution has to be subtracted from the quark distribution,

for each flavour seperately.

Strangeness in the proton is not only observed via the EMC experiment and via the baryon magnetic moments. Independent evidence for strangeness in the proton is given by the Sigma term analysis (5.7), by elastic neutrino proton scattering (5.4), by $s\bar{s}$ meson couplings to nucleons, and by Kaplan's recently discovered strangeness operator (5.25). Unfortunately it is not possible to ask, how much strangeness 'inside' the proton is, because different operators are considered in different experiments. It is also difficult to specify the exact size of the proton. The dispersion relation analysis (5.17) of the πN Sigma term gives the proton a scalar radius (5.18) which exceeds its charge radius by 50%. Interesting, with respect to environment dependent constituent quark masses, is the evaluation of the matrix element $\triangle m_N =< N|c_8 u_8|N >$ (5.11). Usually, the analysis does not respect the nuclear environment by relating this matrix element to the Ξ and Λ mass. Taking instead effective quark masses (4.37) we would find the astonishingly small value $\sigma_N^{\pi(y=0)} = 13 \pm 5 MeV$ (5.15).

A general justification of the NQM approximation in terms of an effective field theory is outstanding. Two models, the chiral constituent quark (CQM) model (5.33) and the Nambu Jona-Lasinio (NJL) model (5.37), are good candidates for the description of hadronic phenomena below the scale where the chiral symmetry is spontaneously broken.

In the NJL model, the quarks are not confined, the gluons being frozen into effective pointlike q-q couplings. The pions which are fundamental entities in the framework of the CQM are considered to be $q\bar{q}$ bound states in the framework of the NJL model.

As one consequence of this distinction we find a different weak coupling in both models: $g_A = 1$ in the CQM and $g_A < 1$ in the NJL model. But because in both models the anomalous magnetic moment of the constituent quark turns out to be negligible, i.e. $\kappa_q = 0$, the equations (4.10) are justified also in the framework of these models. $\kappa_q = 0$ and $g_A = 1$ implies that the constituent quarks can be seen as pointlike elementary degrees of freedom below the scale where the spontaneous chiral symmetry breaking occurs. The almost infinite number of degrees of freedom above this scale, i.e. the partons which make up a constituent quark, are effectively described by a single degree of freedom as soon as the SSB takes place.

Because the corrections to $g_A = 1$ will not be large anyway, we can compare the spin measurements of the EMC with the spin measurements of the magnetic moments. In the first place we observe again that the F/D (or z) value is the same for constituent and current quarks in most quark models and in reality (3.26/4.20). Next we notice that this value, which is independently (but in a mutually consistent way) derived from the β-decays and the magnetic moments of the baryons, is very close to its SU(6) value.

With respect to this characteristic value, SU(6) is not questioned by these experiments. Questionable is, however, the interpretation of the SU(3) symmetric shift of the approximate solution (5.47) from the naively expected solution (5.48). In order to understand such a shift, the process (5.49) in which the surrounding kaon cloud interacts directly with the constituent quark looks most appealing. A similar process may also be responsible for a similar shift in all s-wave groundstate baryons, the decuplet states included. Only the measurement of the spin dependent structure function of the Ω^- can clarify the suspicion that in all SU(3) baryon groundstates the quark spin sum is identical. In this case a general relativistic reasoning for the SU(6) theory becomes feasible, because the internal symmetries commute with the space symmetry.

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