

Polarized intrinsic charm
as a possible solution to the proton spin
problem

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Abstract:

We argue that a large fraction of the proton spin comes from a contribution due to a nonperturbative intrinsic charm component of the proton. An account of this contribution removes an apparent contradiction between the data and exact Kühn-Zakharov low energy theorem. On the other hand, we show that a large intrinsic charm spin component of the proton is implied by recent CLEO data on $B \rightarrow \eta'$ decays. We argue that the proton spin and $B \rightarrow \eta'$ data are manifestations of the same physics with full agreement between two different estimates of the intrinsic charm component of the proton spin. Our results suggest an explanation of the polarized DIS data and have profound consequences for future experiments on charm production off a polarized proton. A microscopic origin of the effect is related to strong well-localized gluon fluctuations (instantons) where the mass of the charmed quark is not a suppression factor at all, as was naively expected earlier.

1 Introduction

The famous proton spin problem is now almost ten years old. A wide interest in the structure of the proton spin was triggered by the 1987 EMC experiment [1] which has measured the first moment of the proton spin structure function $\Gamma_1^p \equiv \int dx g_1^p(x) = 0.126 \pm 0.018$. This result implies that the polarized deep inelastic scattering (DIS) data cannot be explained simply by the valence quark spin components [2], but, in addition, a substantial sea-quark polarization is needed. Such a polarized quark sea seems very miraculous from the point of view of both the relativistic quark model (where valence quarks account for 3/4 of the proton spin) and the parton model (where the helicity conservation prohibits a perturbative sea polarization by hard gluons). This problem, sometimes called “the proton spin crisis”, has stimulated a great deal of interest over the last decade (see e.g. [3] for a review and representative list of references).

Nowadays, it is commonly accepted that these are gluons that play an important role in the spin structure of the nucleon. Depending on a factorization scheme, a gluon contribution to Γ_1^p is either given by an independent anomaly term in the chiral-invariant scheme [4], or included in spin-dependent quark distributions in the Operator Product Expansion (OPE) based gauge-invariant scheme advocated by Jaffe and Manohar [5]. In the Jaffe and Manohar approach, which we will follow, the data are usually interpreted as the manifestation of a negatively polarized s -quark component in the proton. An equivalence of these two descriptions for the first moment Γ_1^p was established by Bodwin and Qiu [6]. While the former presumably makes more contact with one’s intuition and the parton model, the latter is based entirely on QCD and turns out more convenient technically (see the Cheng’s paper in [3] and below).

Irrespective of the question of interpretation of Γ_1^p , the gluon polarization $\Delta G(x)$ in the proton is a very interesting object on its own. This quantity is currently attracting a great deal of attention of both theorists and experimentalists. Among experimental tests on $\Delta G(x)$ a hadronic heavy-quark production (and, in particular, the charm production) is believed to be the best one. The reason for this belief is that within the standard perturbative photon-gluon fusion prescription the charm production amplitude is very sensitive to the gluon distribution in the nucleon. Among future experiments, intended to probe the gluon polarization through this scenario, one should mention the RHIC and HERA- \vec{N} projects on a charm production in polarized $p\bar{p}$ collisions, and the COMPASS experiment at LHC on an open charm production in polarized DIS (see the review by Ramsey in [3] for references and more detail). Existing theoretical expectations for these experiments are all based on the perturbative photon-gluon fusion scenario.

However, while there is little doubt that this mechanism is operative for unpolarized DIS, its relevance may (and must) be questioned in the case of polarized DIS. Indeed, in the latter case one deals with the axial channel instead of the vector one, as it was for usual DIS. On the other hand, it is very well known that physics of axial channels is drastically different from that of vector ones. Unlike the latter, **axial and pseudoscalar channels are strongly influenced by nonperturbative effects**. The best known example of this difference is provided by a nonperturbative breakdown of the Zweig rule in the pseudoscalar nonet. Thus, the perturbative charm production mechanism, valid for usual DIS, certainly cannot be taken for granted in the polarized case. In particular,

strong nonperturbative fluctuations in the axial channel could induce a sizeable “intrinsic charm” fraction of the proton spin, thus substantially changing theoretical predictions for charm production in polarized experiments.

The purpose of this paper is to argue that the latter hypothetical situation is in fact precisely what happens in reality. We will present two independent estimates of the intrinsic charm component of the proton spin, which agree with each other and both indicate that this quantity is quite large. Our main argument is based on recent experimental results, while a second estimate relies on a theoretical result which has been known for a long time, but unfortunately has not received much attention in the literature. As will be shown below, the observation of a large c -quark contribution to the proton spin turns out to be crucial for understanding the polarized DIS data. Our arguments go beyond perturbation theory.

Our main estimate comes from recent data of B-physics (!), which at first sight might seem entirely uncorrelated with the proton spin problem. We mean the new CLEO results on $B \rightarrow \eta'$ decays [7]. Recently, we have shown [8] that η' production in B decays can only be explained by the presence of a large intrinsic charm component $\langle 0 | \bar{c} \gamma_\mu \gamma_5 c | \eta' \rangle$ of the η' . Moreover, a theoretical prediction for this quantity [8] yields a good agreement with the data in both the exclusive [8] and inclusive [9] cases. On the other hand, it is fairly obvious that this property of the η' implies a sizeable nonperturbatively generated intrinsic charm component of the proton spin $\langle p | \bar{c} \gamma_\mu \gamma_5 c | p \rangle$, as soon as there exists a Goldberger-Treiman type relation between these matrix elements. A relation with the B-physics provides us with an estimate of this quantity which, as will be shown below, agrees with an estimate based on a pure theoretical result which we are about to discuss.

Our second estimate is based on a comparison of the DIS data with a low energy theorem due to Kühn and Zakharov (KZ) [10], which is an *exact* (in the chiral limit) formula for the proton matrix element of the topological density. This important result, obtained in 1990, sank into oblivion in subsequent publications. We understand that one of the reasons for this attitude was an apparent sign contradiction between the KZ formula and the data. As we will argue, this contradiction is actually removed by the intrinsic charm component of the proton spin in a very natural way and in full agreement with our previous estimate based on the analysis of B decays. Therefore, the two independent lines of reasoning lead to the same conclusion on a role of the polarized charm in the proton spin.

We thus arrive at a coherent and attractive explanation of the polarized DIS phenomenology. In our opinion, the measured first moment Γ_1^p is determined (we assume the chiral limit) by four cornerstones:

- (1) spontaneous breaking of the $SU_L(3) \times SU_R(3)$ chiral invariance (see below),
- (2) a resolution of the $U(1)$ problem,
- (3) the proton matrix element of the topological density, fixed by the KZ theorem, and
- (4) the intrinsic charm component of the proton spin, which is extracted from data on $B \rightarrow \eta'$ decays.

As we will show below, a consistent treatment of these four ingredients leaves very little (if any) of the mystery of Γ_1^p , and has profound consequences for the future experiments. Our presentation is organized as follows. We start in Sect.2 with a re-interpretation of experimental facts on the polarized DIS. The only difference from the standard analysis is

that we assume a non-zero (but unspecified any further at this stage) value of the matrix element $\langle p|\bar{c}\gamma_\mu\gamma_5c|p\rangle$ which describes the intrinsic charm component of the proton spin. We use the data to select a $SU(3)$ flavor singlet contribution (which is a combination of the anomaly and intrinsic charm terms) to different flavor components of the proton spin. We thus interpret the data as a constraint on this $SU(3)$ singlet contribution. In Sect.3 we present the KZ theorem and then use it to select from this combination a contribution of the intrinsic charm proper. In Sect.4 the same quantity is estimated independently from the data on $B \rightarrow \eta'$ decays. In Sect.5 we briefly discuss consequences of our results for a charm production in polarized experiments. Our conclusions are presented in final Sect.6.

2 $SU(3)$ singlet proton spin from the data

We start with the OPE analysis of the first moment of the polarized g_1 structure function. The only deviation from the standard treatment discussed in detail in [3] is that we keep the charm component of the electromagnetic current of virtual photon. Retaining this term amounts to keeping trace of an intrinsic charm component of the proton spin. In the standard analysis this component is discarded on the grounds that this contribution is expected to appear only at higher orders in α_s from loop effects. As will be shown below, this term is actually $O(\alpha_s^0)$, though it is suppressed by $1/m_c^2$ in agreement with general arguments by Jaffe and Manohar [5]. Therefore, we insist on keeping it in the analysis. On the other hand, at EMC energies $Q^2 \simeq 10 \text{ GeV}^2$, the mass of a struck c -quark can be neglected, and thus the sole effect of taking the charm contribution into account in the OPE analysis¹ is simply reduced to adding a corresponding charm contribution to light u, d, s quark spin components. In the standard nomenclature, we calculate the quantity

$$\Gamma_1^p(Q^2) \equiv \int_0^1 dx g_1(x, Q^2). \quad (1)$$

Using the OPE for a T-product of two electromagnetic currents and selecting an antisymmetric in Lorentz indices contribution, one arrives at the result

$$\Gamma_1^p(Q^2) = \frac{1}{2} \left(\frac{4}{9} \Delta u(Q^2) + \frac{1}{9} \Delta d(Q^2) + \frac{1}{9} \Delta s(Q^2) + \frac{4}{9} \Delta c(Q^2) \right) \left(1 - \frac{\alpha_s(Q^2)}{\pi} + \dots \right). \quad (2)$$

Here Δq (where q is any of the quark flavors) stands for the spin component of the proton due to the quark flavor q

$$s_\mu \Delta q(Q^2) = \langle p|\bar{q}\gamma_\mu\gamma_5q|p\rangle, \quad (3)$$

(s_μ is the proton spin vector) and the operator in (3) is normalized at the normalization point Q^2 . To select a $SU(3)$ singlet contribution to (2), it is convenient to use the combinations

$$g_A^3 = \Delta u - \Delta d, \quad g_A^8 = \Delta u + \Delta d - 2\Delta s, \quad (4)$$

which have no anomalous dimensions and thus can be found from low energy neutron and hyperon beta decays data. Assuming the $SU(3)$ flavor symmetry, the non-singlet

¹Usual unpolarized DIS on intrinsic charm was considered in [11].

couplings are

$$g_A^3 = F + D, \quad g_A^8 = 3F - D. \quad (5)$$

For the values $F = 0.459 \pm 0.008$, $D = 0.798 \pm 0.008$, this yields $g_A^8 = 0.579 \pm 0.025$ [3]. Using relations (4), one can now express the answer for Γ_1^p in terms of the (known) constants g_A^3, g_A^8 and the unknown $SU(3)$ singlet combination

$$\Delta\Sigma(Q^2) \equiv \Delta u(Q^2) + \Delta d(Q^2) + \Delta s(Q^2) \quad (6)$$

plus an intrinsic charm part $\langle p | \bar{c} \gamma_\mu \gamma_5 c | p \rangle$. We obtain

$$\Gamma_1^p = C_{NS}(Q^2) \left(\frac{1}{12} g_A^3 + \frac{1}{36} g_A^8 \right) + \frac{1}{9} C_S(Q^2) (\Sigma(Q^2) + 2\Delta c(Q^2)). \quad (7)$$

As we have said before, the only difference from the standard formulas in Eq.(7) is the presence of an intrinsic charm term Δc . Coefficients C_{NS}, C_S account for perturbative QCD corrections $C_{NS(S)} = 1 - \alpha_s/\pi + \dots$ which can be found in [3].

A global fit to all available data (including the data of SMC, E142 and E143) together with a treatment of higher order QCD corrections, but neglecting the charm contribution gives² [12]

$$\Delta u = 0.83 \pm 0.03, \quad \Delta d = -0.43 \pm 0.03, \quad \Delta s = -0.10 \pm 0.03 \quad (8)$$

with

$$\Delta\Sigma = 0.31 \pm 0.07 \quad (9)$$

at $Q^2 = 10 \text{ GeV}^2$. In our case a difference from this result comes due to retaining the intrinsic charm Δc term in Eq.(7). Therefore, we translate Eq.(9) in the constraint

$$g_A^0 \equiv \Delta\Sigma + 2\Delta c = 0.31 \pm 0.07. \quad (10)$$

(again $Q^2 = 10 \text{ GeV}^2$ is implied). It is convenient at this stage to proceed to derivatives instead of axial currents which stand in (10). Using the result of the OPE for the c -quark bilinear (see Eq.(12) below) and working in the chiral $SU(3)$ limit (which will be implied in what follows), we obtain for $n_f = 3$ light flavors

$$\langle p | n_f \frac{\alpha_s}{4\pi} G \tilde{G} - 2 \cdot \frac{1}{16\pi^2 m_c^2} g^3 G \tilde{G} G | p \rangle = g_A^0 \cdot 2M_p \bar{p} i \gamma_5 p. \quad (11)$$

Here the first term is the standard contribution due to the light quarks. The second term is new, and plays a very important role in what follows. It originates in the c -quark contribution Δc , and technically comes from the derivative of the axial charmed current

$$\begin{aligned} \partial_\mu (\bar{c} \gamma_\mu \gamma_5 c) &= \frac{\alpha_s}{4\pi} G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a + 2m_c \bar{c} i \gamma_5 c \\ &= \frac{\alpha_s}{4\pi} G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a - \frac{\alpha_s}{4\pi} G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a - \frac{1}{16\pi^2 m_c^2} g^3 f^{abc} G_{\mu\nu}^a \tilde{G}_{\nu\alpha}^b G_{\alpha\mu}^c + \dots \end{aligned} \quad (12)$$

²These numbers were slightly updated by Ellis and Karliner in [3]. To our accuracy this difference is unessential.

(see the appendix in [8] for a detailed derivation of this result).

If we were to forget about the second charm-induced term in Eq.(11), we would conclude that the quark contribution to the proton spin is given by the proton matrix element of the topological density (which has to be understood as the limit of a near-forward matrix element for vanishing momentum transfer)³. The corresponding matrix element is fixed by a beautiful result due to Kühn and Zakharov (KZ) [10], which provides us with an *exact* (in the chiral limit) answer for the matrix element of the topological density:

$$\langle p | n_f \frac{\alpha_s}{4\pi} G\tilde{G} | p \rangle = -\frac{2n_f}{3b} 2M_p \bar{p} i \gamma_5 p. \quad (13)$$

As will be shown below, this elegant formula is not sufficient to explain the polarized DIS data. The second charm-induced term in Eq.(11) turns out to be absolutely crucial in this problem. We postpone a corresponding discussion to the next section, while here we want to go one step further and ask the question whether it is possible to find flavor contributions to the proton spin separately. The answer to this question is affirmative. To make contact with previous analyses which have not taken into account the charm contribution, we find it convenient to split it equally between the three light flavor contributions, i.e. to consider “shifted” values $\Delta q' \equiv \Delta q + 2/3 \Delta c$ for $q = u, d, s$. This is in accordance with our definition (10) where Δc simply redefines the singlet contribution $\Delta \Sigma$. Such a procedure amounts to adding an $SU(3)$ singlet piece equally to each of the spin dependent light quark distributions, which was accounted previously by the anomaly term alone without a charmed contribution. We may then translate the result of Ellis-Karliner fit (8) in the values of $\Delta q'$ rather than Δq as it was originally stated in Eq.(8). On the other hand, taking the derivative for each of the three flavor $q = u, d, s$, we obtain

$$\langle p | 2m_q \bar{q} i \gamma_5 q + \tilde{Q} | p \rangle = \Delta q' \cdot 2M_p \bar{p} i \gamma_5 p, \quad (14)$$

where \tilde{Q} stands for the $SU(3)$ singlet part :

$$\tilde{Q} = \frac{\alpha_s}{4\pi} G\tilde{G} - \frac{2}{3} \frac{1}{16\pi^2 m_c^2} g^3 G\tilde{G}G. \quad (15)$$

Here we would like to remind the reader that we work in the chiral limit $m_q = 0$. However, it would be completely erroneous to put $m_q = 0$ in matrix elements (14) from the very

³While, as will be explained below, this neglect is not legitimate, our point here is different. The fact that the light quark fraction of the proton spin $\Delta \Sigma$ is *equivalent* to a particular value of the anomaly matrix element over the nucleon was strongly emphasized by Jaffe and Manohar [5]. This statement is a consequence of the exact equation of motion. A very similar situation holds in another anomalous case: the η' residue $f_{\eta'}$ measures a coupling of the η' to quarks, $\langle 0 | \sum \bar{q}_i \gamma_\mu \gamma_5 q_i | \eta' \rangle = i\sqrt{3} f_{\eta'} q_\mu$. On the other hand, the same parameter gives in the chiral limit a coupling of the η' to gluons: $\langle 0 | n_f \alpha_s / (4\pi) G\tilde{G} | \eta' \rangle = \sqrt{3} f_{\eta'} m_{\eta'}^2$, as a consequence of exact equation of motion. As was stressed in [5], gluon operators cannot appear explicitly in the OPE (2) since there is no gauge invariant gluon operator of spin 1 and twist 2. Thus, in the gauge invariant OPE approach to polarized DIS a gluon contribution is “hidden” (via the subtraction of regulator terms) in quark spin dependent distributions, which therefore do not have a parton interpretation. Later, Bodwin and Qiu have demonstrated [6] that any gauge invariant (i.e. Pauli-Villars) regularization of the photon-gluon scattering diagram automatically yields zero contribution of this diagram to the first moment Γ_1^p , in agreement with the general logic of OPE. Moreover, the approach based on the OPE is more convenient technically than the alternative method [4], since it deals at every step with matrix elements of local gauge invariant operators, which are well defined objects.

beginning. Proceeding this way, we would miss the phenomenon of spontaneous breaking of the chiral $SU(3)_L \times SU(3)_R$ invariance. Goldstone bosons, resulting from this breakdown, have masses $m^2 \sim m_q$, and therefore are the only surviving intermediate states in matrix elements $\langle p|m_q \bar{q}i\gamma_5 q|p \rangle$ in the chiral limit $m_q \rightarrow 0$, which is only taken after a saturation of these matrix elements by goldstones. Technically, it is more convenient not to do it explicitly, but rather use relations (4). Taking the derivatives, we obtain

$$\begin{aligned} \langle p|2m_u \bar{u}i\gamma_5 u - 2m_d \bar{d}i\gamma_5 d|p \rangle &= g_A^3 \cdot 2M_p \bar{p}i\gamma_5 p, \\ \langle p|2m_u \bar{u}i\gamma_5 u + 2m_d \bar{d}i\gamma_5 d - 4m_s \bar{s}i\gamma_5 s|p \rangle &= g_A^8 \cdot 2M_p \bar{p}i\gamma_5 p, \\ \langle p|2m_u \bar{u}i\gamma_5 u + 2m_d \bar{d}i\gamma_5 d + 2m_s \bar{s}i\gamma_5 s|p \rangle &= 0. \end{aligned} \quad (16)$$

We would like to stress that the last of Eqs.(16) implies a resolution of the U(1) problem, i.e. the absence of a $SU(3)$ flavor singlet meson with $m^2 \sim m_q$. We are then able to find exact expressions for the $SU(3)$ flavor dependent parts for each of the light flavor spin distributions separately:

$$\begin{aligned} \langle p|2m_u \bar{u}i\gamma_5 u|p \rangle &= \left(\frac{1}{2}g_A^3 + \frac{1}{6}g_A^8 \right) \cdot 2M_p \bar{p}i\gamma_5 p = (0.725) \cdot 2M_p \bar{p}i\gamma_5 p, \\ \langle p|2m_d \bar{d}i\gamma_5 d|p \rangle &= \left(-\frac{1}{2}g_A^3 + \frac{1}{6}g_A^8 \right) \cdot 2M_p \bar{p}i\gamma_5 p = (-0.532) \cdot 2M_p \bar{p}i\gamma_5 p, \\ \langle p|2m_s \bar{s}i\gamma_5 s|p \rangle &= \left(-\frac{1}{3}g_A^8 \right) \cdot 2M_p \bar{p}i\gamma_5 p = (-0.193) \cdot 2M_p \bar{p}i\gamma_5 p. \end{aligned} \quad (17)$$

In the numerical values in (17) we have implied that the numbers g_A^3, g_A^8 in the chiral limit are not very different from their phenomenological values in the real world (5). We expect an accuracy of this approximation to be rather good⁴. It is curious to note that, despite of their very transparent meaning and exactness in the chiral limit, we were unable to find relations (17) in the literature. Meanwhile, they reveal something very interesting. Indeed, comparing this numbers to Ellis-Karliner fit of experimental results (8) (remember that we have agreed to understand them as “shifted” values $\Delta q' = \Delta q + 2/3\Delta c$), we see that, taken separately, each light flavor spin distribution is given by a large $SU(3)$ dependent part plus a small $SU(3)$ singlet piece whose numerical value is approximately **the same and is very close to 0.1 for all the light flavors**. We thus see that the mystery of Γ_1^p has to a large extent disappeared: the $SU(3)$ flavor dependent parts give the main contributions to each of the light flavor spin distributions Δq , but, on the other hand, these $SU(3)$ variant parts cancel out in the total light flavor contribution $\Delta\Sigma$, as a consequence of resolution of the U(1) problem (see the last of Eqs.(16)).

Therefore, if we managed to establish theoretically the number

$$\frac{1}{2M_p \bar{p}i\gamma_5 p} \langle p| \frac{\alpha_s}{4\pi} G\tilde{G} - \frac{2}{3} \frac{1}{16\pi^2 m_c^2} g^3 G\tilde{G}G|p \rangle \simeq 0.1, \quad (18)$$

(which is simply $1/3 g_A^0$, see Eq.(11)), the Γ_1^p problem would be resolved in QCD terms. Note that constraint (18), resulting from comparison of numbers (17) with Ellis-Karliner

⁴More accurately, it can be shown that a decrease of M_p due to diminishing of the s -quark mass nearly compensates an increase of g_A^3, g_A^8 in the chiral limit. Thus, their product remains approximately the same as in the real world.

fit (8), agrees with Eq.(11) obtained without a specification of different flavor contributions to the quantity $\Sigma + 2\Delta c$. Thus, $SU(3)$ singlet contribution (18) is indeed flavor independent as should be expected. Furthermore, it consists of two parts. Had we knew the anomaly matrix element over the proton independently, we would then be able to extract the intrinsic charm component to the proton spin (the second term in Eq.(18)) from experimental constraint (18). As will be shown in the next section, this knowledge is provided by the KZ theorem, and thus such a separation can indeed be made.

3 Kühn-Zakharov low energy theorem and intrinsic charm in the proton spin

Our aim in this section is to separate the anomaly and intrinsic charm contributions in Eq.(18). To this end, we are going to use a beautiful result by Kühn and Zakharov (KZ) [10] which provides us with an *exact* (in the chiral limit) answer for the proton matrix element of the topological density (the first term in (18)). KZ have shown that the famous dimensional transmutation phenomenon, which lies at heart of QCD low energy theorems, actually also fixes a value of the latter matrix element. In view of the fact that this important result is usually either ignored or doubted in the literature, we would like first to present KZ arguments at length and advocate their correctness. Next, we will use this result to obtain our first estimate of the intrinsic charm contribution to the proton spin.

We start with recalling the well known fact that in a massless asymptotically free theory (such as QCD in the chiral limit) the appearance of a mass scale is related to the so-called dimensional transmutation phenomenon, according to which in this case the only mass parameter in the theory is

$$m \equiv \Lambda \exp\left(-\frac{8\pi^2}{bg^2(\Lambda)}\right), \quad (19)$$

where $b = 11/3N_c - 2/3n_f$ is the first coefficient of the Gell-Mann - Low beta function and Λ stands for a ultraviolet cut-off (N_c is a number of colors). By construction, m does not depend on the cut-off Λ , and all physical parameters (masses, vacuum condensates, etc.) in the chiral limit can only be proportional to a power of the mass parameter (19), as they cannot depend on the cut-off explicitly. Specifying on the proton mass $M_p \sim m$, this means that under the variation of the cut-off

$$\Lambda \rightarrow \Lambda(1 + \varepsilon) \quad (20)$$

a total variation of M_p must vanish:

$$\varepsilon\Lambda \frac{dM_p}{d\Lambda} = 0. \quad (21)$$

On the other hand, the total derivative in respect to Λ can be represented as the sum of a partial derivative and a term corresponding to a variation of the Lagrangian under cut-off shifts (20) :

$$0 = \varepsilon\Lambda \frac{dM_p}{d\Lambda} = \varepsilon\Lambda \frac{\partial M_p}{\partial \Lambda} + \delta_\Lambda M_p, \quad (22)$$

where

$$\delta_\Lambda M_p \bar{p}p = -\langle p | \delta_\Lambda L^{(\Lambda)} | p \rangle = -\langle p | \varepsilon \Lambda \frac{\partial L^{(\Lambda)}}{\partial \Lambda} | p \rangle \quad (23)$$

and $L^{(\Lambda)}$ is a regularized QCD Lagrangian. Its variation under scale transformations (20) is expressed by the conformal anomaly equation through the trace of the energy-momentum tensor $\theta_{\mu\mu}$

$$\varepsilon \Lambda \frac{\partial L^{(\Lambda)}}{\partial \Lambda} = \varepsilon \theta_{\mu\mu} = -\varepsilon \frac{b\alpha_s}{8\pi} G_{\mu\nu}^a G_{\mu\nu}^a. \quad (24)$$

On the other hand, the partial derivative of M_p is fixed by Eq.(19) :

$$\varepsilon \Lambda \frac{\partial M_p}{\partial \Lambda} = \varepsilon M_p. \quad (25)$$

Combining Eqs.(22,23,24,25), we obtain

$$\langle p | -\frac{b\alpha_s}{8\pi} G^2 | p \rangle = M_p \bar{p}p. \quad (26)$$

Eq.(26) is well known and expresses the fact that in the chiral limit the proton mass is given by the conformal anomaly matrix element over the proton. As was pointed out in [10], Eq.(26) can be obtained in a number of ways, thus providing an independent check of the above procedure.

A particularly interesting observation made by KZ [10] was that the very same line of reasoning could be applied in situations when only a fermionic cut-off (specifically the mass M_R of a Pauli-Villars fermion regulator) is varied in two possible ways:

$$M_R \rightarrow M_R(1 + \varepsilon) \quad (27)$$

or

$$M_R \rightarrow M_R(1 + i\varepsilon), \quad (28)$$

where ε is a small and real number. Consider first variation (27). Again, the renormalizability of the theory implies that a total variation of the proton mass under shifts (27) must vanish. On the other hand, a variation of the Lagrangian under a real shift of M_R can be readily found. Indeed, at the one-loop level contributions of fermion and boson regulators to the conformal anomaly are additive, and therefore the variation of the Lagrangian under transformations (27) is

$$\delta_{M_R} L = \varepsilon \cdot \frac{2}{3} n_f \alpha_s G^2, \quad (29)$$

which together with Eq.(26) yields

$$\varepsilon M_R \frac{\partial M_p}{\partial M_R} = -\varepsilon \frac{2n_f}{3b} M_p. \quad (30)$$

Here we come to a most interesting (and sometimes wrongly criticized in the literature) part of the KZ proposal. Consider now the complex regulator mass variation (28). This

variation leads to a γ_5 mass term for the Pauli-Villars fermion R , as Eq.(28) is equivalent to the transformation

$$M_R \bar{R} R \rightarrow M_R \bar{R} R + \varepsilon M_R \bar{R} i \gamma_5 R . \quad (31)$$

Therefore, the variation of the Lagrangian in this case is proportional to the axial anomaly

$$\delta_{iM_R} L = \varepsilon n_f \frac{\alpha_s}{8\pi} G \tilde{G} . \quad (32)$$

Again, a shift in the physical proton mass due to the matrix element of anomaly (32) must be cancelled by a partial derivative in respect to variation (28) (see a comment below). On the other hand, the partial derivative is now determined by Eq.(30) :

$$i\varepsilon M_R \frac{\partial M_p}{\partial M_R} = -i\varepsilon \frac{2n_f}{3b} M_p . \quad (33)$$

Therefore, matrix element of the chiral anomaly (32) is fixed [10] by the requirement of independence of the physical mass of shifts (28) :

$$\langle p | \frac{\alpha_s}{8\pi} G \tilde{G} | p \rangle = -\frac{2}{3b} M_p \bar{p} i \gamma_5 p , \quad (34)$$

that completes the proof of the KZ theorem. As was stressed in [10], the above derivation implies (again) a resolution of the U(1) problem since otherwise the matrix element of interest would not exist in the limit of vanishing momentum transfer because of a massless U(1) goldstone contribution. Moreover, while there may exist higher order corrections to Eq.(34), this result is exact in the chiral limit at least to the one-loop accuracy. Its validity was checked by KZ in the case of supersymmetric QCD, and may also be tested in solvable models.

As the negative sign in Eq.(34) is so important for what follows, it is instructive to have an alternative derivation of this result, where the sign would be explained in a simple and intuitive way. We now present such a derivation.

The physical meaning of Eq.(26) is very simple: it says that in the chiral limit a considerable part of the nucleon mass $\frac{11N_c}{3b} M_p$ comes from gluons, while the term related to quarks (more precisely, to their regulator fields) is equal to $\frac{-2n_f}{3b} M_p$. Let us introduce chiral combinations of the regulator fields in the standard way:

$$R_l = \frac{1}{2}(1 + \gamma_5)R, \quad R_r = \frac{1}{2}(1 - \gamma_5)R. \quad (35)$$

Transformation properties of the R_l , R_r fields under chiral rotations $\sim \exp(i\alpha\gamma_5)$ have the very simple form:

$$\bar{R}_r R_l \rightarrow \exp(i\alpha) \bar{R}_r R_l, \quad \bar{R}_l R_r \rightarrow \exp(-i\alpha) \bar{R}_l R_r \quad (36)$$

The requirement of reparametrization invariance under chiral rotations (36) in the limit $m_q = 0$ can be expressed by equations

$$\langle p | -M_R \bar{R}_r R_l e^{i\alpha} | p \rangle = M_p^{(f)} e^{i\alpha} \bar{p}_r p_l, \quad \langle p | -M_R \bar{R}_l R_r e^{-i\alpha} | p \rangle = M_p^{(f)} e^{-i\alpha} \bar{p}_l p_r, \quad (37)$$

which should be valid for arbitrary α . In this formula $M_p^{(f)}$ is a part of the nucleon mass which comes from the fermion regulator field⁵: $M_p^{(f)} = (\frac{-2}{3b}) M_p$. One can easily check this formula once again using an expansion for the scalar density at $M_R \rightarrow \infty$:

$$\langle p | -M_R \bar{R} R | p \rangle \rightarrow \langle p | \frac{\alpha_s}{12\pi} G^2 + 0(1/M_R^2) | p \rangle = M_p^{(f)} \bar{p} p = \frac{-2}{3b} M_p \bar{p} p, \quad (38)$$

which is exactly the contribution of one flavor to complete expression (26). A similar calculation for the pseudoscalar part of Eq.(37) gives the result (for convenience, we multiply both parts of the equation by the factor i):

$$\langle p | -M_R \bar{R} i \gamma_5 R | p \rangle \rightarrow \langle p | \frac{\alpha_s}{8\pi} G \tilde{G} + 0(1/M_R^2) | p \rangle = M_p^{(f)} \bar{p} i \gamma_5 p = \frac{-2}{3b} M_p \bar{p} i \gamma_5 p, \quad (39)$$

where we used the standard expansion $-M_R \bar{R} i \gamma_5 R \rightarrow \frac{\alpha_s}{8\pi} G \tilde{G}$ for the large mass M_R , see Eq.(12). Expression (39) is precisely KZ theorem (34). Now we understand the sign in KZ formulae (34), (39) very well: it has the **minus** sign because a fermion field gives a negative contribution to the nucleon mass in comparison with the main term originating from gluons. By the same reasons this term is suppressed in the large N_c limit. This statement is very clear and unambiguous.

We would now like to guess a reason why the KZ formula is mostly ignored in the literature. The reason is the negative sign of matrix element (34), which is absolutely crucial in what follows. Had we forgotten about the second term in Eq.(18), we would conclude that result (34) contradicts the data as it is of a different sign. Our answer to this objection is that it is the second charm-induced term in Eq.(18) that makes the whole expression positive. As will be shown in the next section, an independent estimate of this term, based on results on $B \rightarrow \eta'$ decays, confirms this claim⁶.

Finally, we are ready to use the KZ theorem given by Eq.(34) to find a charm-induced contribution to the proton spin from experimental constraint (18). We obtain

$$\frac{1}{2M_p \bar{p} i \gamma_5 p} \langle p | -\frac{1}{16\pi^2 m_c^2} g^3 G \tilde{G} G | p \rangle \simeq 0.3 \quad (40)$$

We thus conclude that experimental result (18), taken together with KZ formula (34), requires a substantial contribution of intrinsic charm to the proton spin, which is about twice larger than the total **singlet** light flavor contribution given by the first term in Eq.(18). We should note that, taken separately, the u and d quark contributions (17) are numerically considerably larger than the c -quark term. It is the singlet combination which is smaller than the charmed quark contribution. Such a smallness of the singlet combination seems natural within the large N_c approach: it should vanish according

⁵Different n_f fermions multiply this result by the factor n_f .

⁶Here we would like to mention that in alternative approach to the anomaly matrix element over the proton, which is based on U(1) Goldberger-Treiman relations within a "two-component" approach of Ref.[13] (see also references in [3]), a sign of matrix element (34) comes as a result of interplay between the η' and a composite gluon ghost field interactions with the nucleon, which may in principle lead to any signature. Yet, in the model-independent KZ method this sign is fixed to be negative. A discussion of possible consequences of this observation for the formalism of Ref.[13] would lead us far beyond the scope of this paper.

to KZ formula (34). We consider this as a reasonable qualitative explanation of our result: the charm contribution becomes competitive in comparison with the $SU(3)$ singlet light quark term because the latter is simply suppressed by the factor $1/N_c$ on a natural scale. Nevertheless, we understand that this result may be not easy to reconcile with one's intuition. On the other hand, the first term in Eq.(18) is the matrix element of a total derivative, i.e. a very unusual object. The fact that it does not vanish can be understood as resulting from interactions in a gauge-variant sector of the theory. The physical intuition can hardly work in such untypical situation. Thus, to support the above conclusion on a role of intrinsic charm in the proton spin, we need to find matrix element (40) in an independent way. If number (40) was established without a reference to the polarized DIS data, the latter would be completely explained in terms of theoretically calculated matrix elements (34) and (40) (see Eq.(18) above). We will now argue that estimate (40) can be confirmed within a different line or reasoning, based on a very different experiment, which is completely independent of both constraint (18) and KZ result (34).

4 Intrinsic charm in the proton and $B \rightarrow \eta'$ data

We now proceed to an independent estimate of matrix element (40) which measures a contribution of the intrinsic charm component of the proton spin to the first moment $\int dx g_1(x)$ of the polarized structure function $g_1(x)$. Somewhat unexpectedly, such an estimate can be obtained from new data of B-physics. As will be argued in this section, experimental information on $B \rightarrow \eta'$ decays can be used to make conclusions on a role of the intrinsic charm component $\langle p | \bar{c} \gamma_\mu \gamma_5 c | p \rangle$ in the proton spin. We will show that this quantity may be related to the intrinsic charm component $\langle 0 | \bar{c} \gamma_\mu \gamma_5 c | \eta' \rangle$ of the η' . The latter matrix element is actually probed experimentally as will be argued below. Thus, in what follows we will first spend some time discussing the B-physics, and then demonstrate how this can be helpful in the problem of interest.

Recently, CLEO collaboration has reported [7] results of measurements of inclusive and exclusive production of the η' in B-decays :

$$Br(B \rightarrow \eta' + X ; 2.2 \text{ GeV} < E_{\eta'} < 2.7 \text{ GeV}) = (7.5 \pm 1.5 \pm 1.1) \cdot 10^{-4} , \quad (41)$$

$$Br(B \rightarrow \eta' + K) = (7.8_{-2.2}^{+2.7} \pm 1.0) \cdot 10^{-5} . \quad (42)$$

Here the inclusive branching ratio contains the acceptance cut intended to reduce a background from events with charmed meson interactions in a final state. At first sight, the above numbers might seem quite innocent. However, simple calculations [8, 9] reveal that these data are in severe contradiction with a standard view of the process at the quark level as a decay of the b -quark into light quarks, which could be naively suggested keeping in mind the standard picture of η' as a $SU(3)$ singlet meson made of the u -, d - and s -quarks. In this picture the $B \rightarrow \eta'$ amplitude must be proportional to the Cabbibo suppression factor V_{ub} and, as a result, the standard mechanism of the $B \rightarrow \eta'$ transition yields numbers which are by two orders of magnitude (!) smaller than the data for both the inclusive and exclusive cases. Thus, there must be something beyond this standard picture. Furthermore, inclusive decays are usually dominated by few-particle

final states. Therefore, any meaningful mechanism, intended to explain abnormally large numbers (41,42), should deal with both the inclusive and exclusive data at the same time.

A particular mechanism, suggesting a unified description of both the inclusive and exclusive modes, was proposed in our recent papers [8, 9]. We have shown that at the quark level $B \rightarrow \eta'$ decays can be described by the Cabbibo favored $b \rightarrow c\bar{c}s$ process followed by a transition of $\bar{c}c$ into the η' . The latter transition is possible due to an intrinsic charm component of the η' , and quantitatively can be expressed through the matrix element

$$\langle 0 | \bar{c} \gamma_\mu \gamma_5 c | \eta'(q) \rangle \equiv i f_{\eta'}^{(c)} q_\mu. \quad (43)$$

A reason why this matrix element may be non-zero can be explained as follows. As the c -quark is heavier than the η' , it may only exist in the η' in a loop. Such a loop with the heavy c -quark can be evaluated in terms of gluon fields, populating the η' , by using the background field technique, see Eq.(12). Therefore, matrix element (43) may not vanish due to $\bar{c}c \leftrightarrow \text{gluons}$ transitions.

Prior to a theoretical calculation of the charm current residue $f_{\eta'}^{(c)}$ into the η' , it is instructive to find this quantity “experimentally”. By this we mean a value of $f_{\eta'}^{(c)}$ needed for the above mechanism to explain the data. Simple calculations show that the number [8, 9]

$$f_{\eta'}^{(c)} \simeq 140 \text{ MeV} \quad (\text{“exp”}) \quad (44)$$

yields a good agreement with the data in both the inclusive and exclusive modes. On the other hand, this value may look uncomfortably large as it is only a few times smaller than the analogously normalized residue $\langle 0 | \bar{c} \gamma_\mu \gamma_5 c | \eta_c(q) \rangle = i f_{\eta_c} q_\mu$ with $f_{\eta_c} \simeq 400 \text{ MeV}$ known experimentally from the $\eta_c \rightarrow \gamma\gamma$ decay. Intuitively, these numbers could be expected to differ much more drastically as, in contrast to f_{η_c} , the residue $f_{\eta'}^{(c)}$ is a double suppressed amplitude. It is Zweig rule-violating and besides contains the $1/m_c^2$ suppression factor, as it comes from loop effects. Therefore, one could expect it to be very small. In reality it is not small. There are two reasons for this. First, m_c is not very large on the hadronic 1 GeV scale. Second, and more important, the Zweig rule itself is badly broken down in vacuum 0^\pm channels. Of course, such a breakdown contradicts a naive large N_c counting where a non-diagonal transitions should be suppressed in comparison with diagonal ones. However, a more careful analysis [14, 8] reveals that the large N_c picture and breakdown of the Zweig rule in fact peacefully co-exist: while the large N_c description is quite accurate for the η' , an extent to which the Zweig rule is violated in η' yields a large residue $f_{\eta'}^{(c)}$. We stress that the phenomenon of the breakdown of the Zweig rule in vacuum 0^\pm channels is well known and understood [14], and many phenomenological examples of corresponding physics have been discussed in the literature, see e.g. [14, 15]. The residue $f_{\eta'}^{(c)}$ (which is fundamentally important for our estimates) is another manifestation of the same physics.

We now proceed to a theoretical analysis of the residue $f_{\eta'}^{(c)}$ defined by Eq.(43). Unfortunately, a detailed theoretical consideration of this quantity would require a repetition of our original paper [8] almost at full length. We thus refer to Ref.[8] for details, while here we would like to give an idea and flavor of our method. It is convenient to start with

taking the derivative in Eq.(43). Then we obtain

$$f_{\eta'}^{(c)} = \frac{1}{m_{\eta'}^2} \langle 0 | 2m_c \bar{c} i \gamma_5 c + \frac{\alpha_s}{4\pi} G_{\mu\nu} \tilde{G}_{\mu\nu} | \eta' \rangle. \quad (45)$$

Since the c -quark is heavy, one can use the Operator Product Expansion in inverse powers of the c -quark mass (the heavy quark expansion)

$$2m_c \bar{c} i \gamma_5 c = -\frac{\alpha_s}{4\pi} G_{\mu\nu} \tilde{G}_{\mu\nu} - \frac{1}{16\pi^2 m_c^2} g^3 f^{abc} G_{\mu\nu}^a \tilde{G}_{\nu\alpha}^b G_{\alpha\mu}^c + \dots \quad (46)$$

(see the appendix in [8] for a detailed derivation of this result). Further terms in expansion (46) are neglected in what follows (see however below). We have thus reduced the problem to a calculation of the matrix element of the purely gluonic operator:

$$f_{\eta'}^{(c)} = -\frac{1}{16\pi^2 m_{\eta'}^2} \frac{1}{m_c^2} \langle 0 | g^3 f^{abc} G_{\mu\nu}^a \tilde{G}_{\nu\alpha}^b G_{\alpha\mu}^c | \eta' \rangle. \quad (47)$$

It may seem at first sight that Eq.(47) is of little help since matrix elements of gluon operators are difficult to calculate. A situation with the η' is, however, exceptional, as the η' is strongly coupled to gluons and to some extent can be viewed as a remnant of imaginary purely gluonic world. In effect, matrix element (47) is amenable to a theoretical study which essentially reduces to a nonperturbative analysis of pure Yang-Mills theory. Closely following old ideas due to Witten[16] and Veneziano[17] and using some additional (nonperturbative) arguments, we have managed to estimate matrix element (47) [8]:

$$f_{\eta'}^{(c)} \simeq \frac{3}{4\pi^2 b} \frac{1}{m_c^2} \frac{\langle g^3 G^3 \rangle_{YM}}{\langle 0 | \frac{\alpha_s}{4\pi} G_{\mu\nu} \tilde{G}_{\mu\nu} | \eta' \rangle}. \quad (48)$$

Therefore, we have related the residue of the charmed axial current into the η' with apparently completely unrelated quantity which is a cubic gluon condensate in the pure Yang-Mills theory (we notice that the matrix element of topological density which appears in (48) is known $\langle 0 | (\alpha_s/4\pi) G\tilde{G} | \eta' \rangle \simeq 0.04 \text{ GeV}^3$ [14]). Using all currently available information regarding the vacuum condensate $\langle g^3 G^3 \rangle_{YM}$ in gluodynamics, we have arrived at the numerical estimate [8]

$$f_{\eta'}^{(c)} = (50 - 180) \text{ MeV}. \quad (49)$$

In spite of a large uncertainty of this result, its main source is nevertheless well localized and related to a poor knowledge of the cubic condensate in pure gluodynamics. One can see that our theoretical estimate (49) agrees within errors with “experimental” value (44).

By reasons mentioned above, it is extremely important to calculate matrix element (47) in an independent way within a QCD-based model in order to test the general idea. Such a calculation based on the instanton-liquid model (see e.g. [18] for a review and references to original papers) has been carried out [19] with the following preliminary, but promising result: the matrix element of the three-gluon operator as defined by Eq.(47) is very close to “experimental” value (44). We should remind the reader that this model is extremely successful in description of low-energy hadronic properties, and a typical

accuracy of this approach is not worse than 30 %. A qualitative, microscopic explanation of the enhancement of matrix element (47) is the following. The model suggests the existence of relatively small strongly interacting instantons with a strong gluon field inside the instanton. This field is so strong that a relevant parameter gG/m_c^2 is not small numerically as one could naively expect, but rather is of order one. We are quite confident in calculations [19] because a similar calculation with the standard gluon current $\tilde{G}G$ is in a good agreement with the existing phenomenological data on the η' [18].

We interpret all these results as the evidence that the suggested mechanism indeed explains the data on $B \rightarrow \eta'$ given by numbers (41,42). Therefore, in what follows we consider the intrinsic charm component of the η' , described by Eqs.(43,44,48,49), as an established property of the η' which is probed in experiments and, on the other hand, is understood theoretically. Still, in view of a poor accuracy of our theoretical prediction (49), we will use “experimental” number (44) in our subsequent estimates.

We are now in a position to make our second estimate of an intrinsic charm component of the proton spin. Consider the matrix element ($q = p' - p$, other notations here are self-explanatory)

$$\langle N(p') | \bar{c} \gamma_\mu \gamma_5 c | N(p) \rangle = \bar{u}(p') \left[\gamma_\mu \gamma_5 h_1(q^2) + q_\mu \gamma_5 h_2(q^2) \right] u(p) . \quad (50)$$

Taking the derivative and using Eq.(46), we obtain

$$\langle N(p') | - \frac{1}{16\pi^2 m_c^2} g^3 G \tilde{G} G | N(p) \rangle = i \bar{u}(p') \gamma_5 u(p) \left[2M_N h_1(q^2) + q^2 h_2(q^2) \right] . \quad (51)$$

Assuming now a η' dominance in this matrix element⁷ and using the absence of pole singularities in the second form factor h_2 , which ensures a resolution of the U(1) problem, we arrive at the Goldberger-Treiman type relation

$$h_1(0) = \frac{1}{2M_N} g_{\eta'NN} f_{\eta'}^{(c)} \quad (52)$$

or, equivalently,

$$\frac{1}{2M_N \bar{u}(p) i \gamma_5 u(p)} \langle N(p) | - \frac{1}{16\pi^2 m_c^2} g^3 G \tilde{G} G | N(p) \rangle = \frac{1}{2M_N} g_{\eta'NN} f_{\eta'}^{(c)} . \quad (53)$$

Note that owing to Eq.(46) the anomaly term cancels out in (53)⁸. Unfortunately, the precise value of $g_{\eta'NN}$ is not known, and phenomenological estimates of the coupling $g_{\eta'NN}$

⁷Note that such a saturation becomes exact in the large N_c limit.

⁸This is in contrast to the standard case with U(1) Goldberger-Treiman relations for the $SU(3)$ singlet light flavor current discussed in Ref.[13]. The anomaly term is a total derivative, and the fact that it gives a non-zero contribution to the proton matrix element is due to the presence of non-pole subtraction terms in dispersion relations, which have the same nature as subtraction terms in a correlation function of the topological density [17], and can be assigned to a ghost contribution. This is the reason why the proton spin has a “two-component” form in the approach of Ref.[13]. As a result, the light quark component of the proton spin is expressed in terms of two unknown quantities, neither of which is directly measurable. For these reasons we prefer to use the exact KZ theorem for the anomaly matrix element rather than results from [13]. On the contrary to the case [13], the anomaly term does not show up in Eq.(53), and this is why it is the physical coupling constant $g_{\eta'NN}$ that stands there.

vary from $g_{\eta'NN} \simeq 3$ to $g_{\eta'NN} \simeq 7$ [20]. Assuming the same range for this quantity in the chiral limit and using numerical value (44), we arrive at the estimate

$$\frac{1}{2M_N \bar{u}(p) i\gamma_5 u(p)} \langle N(p) | - \frac{1}{16\pi^2 m_c^2} g^3 G \tilde{G} G | N(p) \rangle = 0.2 - 0.5. \quad (54)$$

This result has a large uncertainty mainly due to a poor knowledge of $g_{\eta'NN}$. In view of a large uncertainty in $g_{\eta'NN}$, we have neglected a perturbative evolution of $f_{\eta'}^{(c)}$.

A few remarks are in order. First, as we already mentioned, the microscopic picture suggests that the matrix element under consideration is not small because of gluon fluctuations which could be large. In different words, the parameter of the heavy quark expansion gG/m_c^2 could be of order one. Such an explanation might imply a bad convergence of expansion (46) where we keep the first term and neglect all the rest. A scientific answer to the question of convergence of series (46) would require the knowledge of higher dimensional matrix elements which are not available at the moment. Alternatively, one could try to address this issue within the instanton liquid approach [18]. However, irrespectively to an outcome of such an analysis, we could consider “experimental” value of $f_{\eta'}^{(c)}$ (44), which effectively takes into account all powers of the $1/m_c$ expansion. Therefore, all of them are also effectively included in the left hand side of Eq.(54). Thus, a possible bad convergence of expansion (46) does not affect our results at all. Hence we see that our final estimate (54) confirms the above conclusion on a large c -quark component in the proton spin, as its sign and order of magnitude agree with phenomenological constraint (40). The large magnitude of the charmed current residue $f_{\eta'}^{(c)}$ into the η' is the main factor giving rise to large number (54). We emphasize again that estimate (54) is obtained in an entirely independent of the DIS phenomenology and KZ theorem way, and relies on the data and theoretical results of B-physics. If we fix the $\eta'NN$ coupling constant in the chiral limit by “experimental” constraint (40), it is consistent within uncertainties with independent theoretical prediction (54). We thus conclude that the polarized DIS data agree within errors with theoretical results expressed by Eqs.(18), (40) and (54). This implies that the main contribution to Γ_1^p is due to the charm component of the proton spin. We now proceed to a discussion of possible ways to test this conclusion.

5 Signals for intrinsic charm in polarized experiments

In this section we would like to argue that our explanation of the polarized DIS data has very definite consequences for future polarized experiments which thus will be able to test the picture suggested by this paper. We will only deal with qualitative aspects of expected phenomena, which follow immediately from our results. A more detailed analysis is beyond the scope of this paper and deserves a separate study.

We have found that the light quark and charm contributions to the first moment Γ_1^p have different signs and magnitudes. Here and in what follows any references to the proton spin are understood in the sense of the singlet part of the spin operator $\Delta\Sigma + 2\Delta c$ only. While the former is negative, the latter is positive with about a twice larger magnitude. The immediate consequence of this fact is that an integral asymmetry for events with no charm in the final state is expected to be negative with approximately the same magnitude

as a total integral asymmetry. This prediction may hopefully be tested experimentally, provided the two types of final states can be distinguished.

Yet, a most interesting test of our scenario is related to a charm distribution in charmed final states. We find that the charm contribution to Γ_1^p is about an order of magnitude larger than results obtained within perturbative schemes [21]. The difference from previous analyses comes due to strong nonperturbative fluctuations in the axial $SU(3)$ singlet channel, which lead to a large magnitude of the intrinsic charm component of the proton spin. Our definition of the latter quantity matches the well known definition [22] of the intrinsic charm in terms of a Fock state description as a multiple connected loop of the c -quark in the proton. In our approach this property is used in a constructive way. Furthermore, one can expect that the intrinsic charm component of the proton spin will be traced in polarized experiments. We note that in the case of unpolarized DIS various mechanisms, which are able to identify the intrinsic charm in the proton, have been discussed in the literature. In particular, they include a liberation of the charm as a result of scattering on light quarks [23], or direct DIS on the intrinsic charm [11]. The intrinsic charm contribution to the proton spin, considered in this paper, corresponds to the latter case.

There exists a clear distinction between pictures of a charm distribution in the final state corresponding to different mechanisms of the charm production in DIS. A vast literature [24] is devoted to the charm production in polarized experiments as an effective tool to measure the gluon spin dependent distribution $\Delta G(x, Q^2)$. This proposal is based on the perturbative photon-gluon fusion (PGF) mechanism which assigns a production of the $\bar{c}c$ pair to the hard subprocess $\gamma^*g \rightarrow \bar{c}c$. It is undoubtedly true that the PGF process adequately describes a heavy quark production in the case of usual unpolarized DIS, and expresses it in terms of the unpolarized gluon distribution. Moreover, in the unpolarized case the manifestation of the intrinsic charm component of the nucleon is still an open problem [11, 23]. We believe that the underlying reason for a relatively subdominant role of the intrinsic charm in usual DIS is a strong suppression of nonperturbative effects in vector channels. However, as we argued in Introduction, it would be potentially dangerous to automatically transfer concepts and results, valid for unpolarized DIS, to the polarized case. Although at the perturbative level results look similar, these arguments ignore important differences between the two cases, which only show up beyond perturbation theory. That this way of thinking is fallacious can be most clearly illustrated on the example of the pseudoscalar nonet. The Zweig rule is 100 % violated there by nonperturbative effects, though no difference from the Zweig rule conserving case of vector mesons is seen in perturbation theory.

As, in contrast to the perturbative PGF mechanism, polarized DIS on the intrinsic charm does not have a strong α_s suppression, we expect that the charm production in polarized experiments will be overwhelmed by events due to intrinsic charm mechanisms. Quantitative predictions in this case are troublesome as, in addition to direct DIS on the intrinsic charm, there exist other possibilities for the extrinsic and intrinsic charm production in DIS [23]. Model independent methods needed for their evaluation are not available at present, though estimates with model light cone wave functions for the intrinsic charm [22] are possible. Nevertheless, different mechanisms for the charm production can be disentangled on the basis of an event topology analysis [11, 23]. The PGF mechanism

corresponds to two jets (plus a spectator) events with a charmed hadron (typically a D -meson) carrying a large fraction of the photon energy in the current jet and an anticharm (e.g., \bar{D}) hadron in the target jet. A topology of these PGF events is expected to deviate from the planar one [11]. The $\bar{c}c$ pair has a low invariant mass as the fractional energy of the proton carried by the gluon is typically small. The transverse momentum and invariant mass of the pair will grow with the photon virtuality Q^2 . As has been argued above, we expect that these events will be parametrically suppressed by powers of α_s in comparison to events reviving the intrinsic charm in the nucleon.

As for DIS on the intrinsic charm, one can distinguish between two types of processes. The first one is an indirect process when the photon scatters on a light quark with a subsequent liberation of the charm. In this case both the charm quark and antiquark are in the target fragmentation region, and can hadronize into both open charm (e.g. D, Λ_c) or hidden charm (J/ψ , etc.) hadrons. In this process, the intrinsic charm shows up through a nonperturbative final state interaction. Charmonium states, produced in this region, will presumably have a large fraction of the proton momentum with $p_\perp \sim m_c$ and independent of Q^2 . On general grounds, one may expect an excessive production of S-wave quarkonia (e.g. η_c) in comparison to P-wave states.

The second process is direct DIS on the intrinsic charm quark [11]. The target charm quark is emitted at large forward rapidities, while the scattered quark has somewhat smaller near forward rapidities. The difference in rapidities grows with Q^2 . The average transverse momentum of the $\bar{c}c$ pair is not expected to grow with Q^2 . Thus, our results allow one to suggest that a substantial (presumably dominant) part of charm events in polarized experiments at large Q^2 will be an open charm hadrons (or S-wave quarkonia) production with low $p_\perp \sim m_c$.

Finally, we note that all said above on the charm production can be extended to beauty production. The latter is suppressed relatively to the former by the factor $m_b^2/m_c^2 \simeq 10$, which may imply that the beauty can be produced nonperturbatively via the intrinsic beauty mechanism at approximately the same level as the charm production through the perturbative PGF process.

6 Conclusions

In this paper we have presented a somewhat unorthodox point of view of the data on the first moment $\int dx g_1(x)$ of the spin $g_1(x)$ structure function measured in polarized DIS experiments. We have abandoned attempts to explain the data within the quark model or related to it field theoretical models, and reformulated the problem in terms of particular matrix elements of quark and gluon operators. In our opinion, none of the phenomena, which are of crucial importance for understanding the proton spin problem (spontaneous breaking of the chiral invariance, a resolution of the U(1) problem, the dimensional transmutation, etc.), can be adequately described within the parton or quark model. On the contrary, all of them are essentially nonperturbative, and need to be addressed within nonperturbative QCD. We have further shown that the data can be interpreted as a constraint on the sum of matrix elements of the anomaly and intrinsic charm operators. The latter object is a new principal element of our analysis, which has

not been considered hitherto in the literature in the content of the proton spin problem. Still, we have argued that the intrinsic charm component of the proton spin is quite large and constitutes a main contribution to the first moment Γ_1^p . Its large magnitude is related to strong nonperturbative fluctuations in the axial $SU(3)$ singlet channel. Moreover, the inclusion of the intrinsic charm proton spin in the analysis is absolutely necessary as it is the only way to reconcile the polarized DIS data with both an exact Kühn-Zakharov low energy theorem on one side, and new data and theoretical results on $B \rightarrow \eta'$ decays on the other. We believe that these seemingly uncorrelated problems are actually tightly connected. The situation reminds us the J/ψ discovery in 1974, when a charmonium state (“hidden charm”) was observed simultaneously in e^+e^- collisions at SLAC [25] and at the proton machine at Brookhaven [26]. We believe that we are now facing a similar case, when different experimental groups see the “intrinsic charm” in polarized DIS (10) and B-decays (41,42) simultaneously.

We feel that at a microscopic level, the new effect of large intrinsic charm fluctuations originates in the instanton physics [18], where strong gluon fields in the singlet channel are able to lead to very unexpected phenomena. Our results seem to imply a kind of universality of this nonperturbative physics with universality classes defined by quantum numbers relevant for a process considered, but not concrete particles involved.

Our explanation of the polarized DIS data has very definite consequences for the future polarized experiments, and thus can be tested there.

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