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Peripheral spin structure and axial couplings of baryons

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We apply the relativistic light-cone perturbation theory to a flow of the helicity of baryons to the angular momentum of the mesonic cloud. Starting with the $SU(6)$ symmetry, we find good description of the axial couplings in β-decay of hyperons. In the light-cone technique, the pattern of the spin-flip of nucleons differs significantly from that expected in the conventional non-relativistic models.

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1. The subject of this paper is a flow of the nucleon helicity from the constituent quarks to the angular momentum of the mesonic cloud. In the charge symmetric γ_5-theory the emission of pseudo-scalar mesons flips the nucleon spin. Then one expects that a part (if not all) of the proton spin is carried by the angular momentum of the P-wave π, K or η meson. The issue has become topical in view of the EMC spin crisis [1]. One of the major sources of information on the spin structure of nucleons is the deep inelastic scattering (DIS) of polarized muons on the polarized nucleons. The quantity measured is the the spin-dependent proton structure function, $g_1^p(x, Q^2)$. g_1^p is related to the proton axial-current matrix element as [2]:

$$2m_N s_\mu \cdot \int_0^1 dx\, g_1^p(x, Q^2) = \frac{1}{2}\left\langle p, s \middle| \sum_q e_q^2 \bar{q}\gamma_\mu\gamma_5 q \middle| p, s \right\rangle. \tag{1}$$

Here e_q^2 is the charge squared of a quark flavour q.

The contribution Δq of a given quark flavour to the spin of the proton is given by the forward matrix element

$$\langle p, s|\bar{q}\gamma_\mu\gamma_5 q|p, s\rangle = 2m_p s_\mu \Delta q, \tag{2}$$

where $|p, s>$ denotes a proton of momentum p, spin s with $s^2 = -1$, $s \cdot p = 0$, normalized to $\langle p, s|p', s\rangle = 2E(2\pi)^3\delta^3(\vec{p} - \vec{p}')$. Then what we need is the flavour singlet combination $\Delta u + \Delta d + \Delta s = g_A^0$ which in the simple constituent quark model is interpreted as the fraction of the proton's spin carried by quarks.

TGV analysis [3] implies that the axial anomaly prevents from attributing the singlet axial charge only to quarks. Discussion of the anomaly effects goes beyond the scope of this paper [4] . We concentrate on the classicall part of the problem: what is a possible axial charge coming from the constituent quarks and what is a role of chiral corrections in the flow of nucleon's spin from the baryonic core of nucleons to the pionic cloud.

Mesonic corrections to the β-decay coupling constants were first calculated in [5]. The lowest-order process (Fig.1) involving a single π^0 has been considered to explain the observed departure of the ratio of the Gamov-Teller and Fermi couplings from unity. Chew's massive fermion theory, involving a cutoff, has been exploited. In the more recent papers on the subject [6] the non-relativistic Cloudy Bag Model is discussed. In the chiral perturbation theory [7] baryons also appear as heavy static fields.

In this paper we study the renormarization of axial couplings of baryons making use of the consistent relativistic light-cone Fock-state approach. We study $SU(6)$ and $SU(3)$ symmetry breaking effects, which with the phenomenologically determined parameters of the light-cone wave functions correspond to $SU(3)$ symmetry breaking by the mass splitting of baryons and mesons. We find that in the light-cone approach the spin-flip pattern differs from that in the simple non-relativistic model, the difference being due to the Fock components in which mesons carry large fraction of baryon's momentum. One interesting observation is that we find good agreement with the hyperon decay axial couplings starting from the classic $SU(6)$ summetry.

2. Ellis and Jaffe [8] noticed that $SU(3)$ octet combinations of Δu, Δd and Δs which are free of anomalous divergences can be related to measured rates of Gamov-Teller transitions between octet baryons. Then, in terms of $SU(3)$ invariant symmetric and antisymmetric parameters, D and F,

$$2\left\langle p, s|A_\mu^3|p, s\right\rangle = 2m_p s_\mu g_A^3 = 2m_p s_\mu(F + D) \tag{3}$$

$$2\sqrt{3}\left\langle p, s|A_\mu^8|p, s\right\rangle = 2m_p s_\mu g_A^8 = 2m_p s_\mu(3F - D), \tag{4}$$

where g_A^3 and g_A^8 are the nucleon axial-vector charges. By definition, $A_\mu^a = \bar{q}\gamma_\mu\gamma_5 T^a q$, and $tr\left(T^a T^b\right) = \frac{1}{2}\delta^{ab}$. The Ellis-Jaffe ansatz, $\Delta s = 0$, would predict $g_A^0 = g_A^8$.

Consider first the proton axial charge g_A^8 renormalization within the pseudoscalar γ_5-theory involving baryon octet and 0^- meson octet. Parametrizing the matrix element of the renormalized current in the form

$$2\sqrt{3}\left\langle p|A_\mu^8|p\right\rangle = \frac{1}{Z_N}\left(\alpha^8 + \beta_\pi^8\lambda_\pi + \beta_K^8\lambda_K + \beta_\eta^8\lambda_\eta\right)\bar{u}_p\gamma_\mu\gamma_5 u_p, \tag{5}$$

we find: $\alpha^8 = 3F - D$, $\beta_\pi^8 = 3(3F - D)$, $\beta_K^8 = -(16/9)D$ and $\beta_\eta^8 = 3F - D$. The first term in (5) is the tree level proton isosinglet axial coupling. The second one presents contribution of the πN Fock component of the proton (Fig.1a).

In the chiral perturbation theory the physical nucleon contains the baryonic core and baryonic core + meson(s) Fock states. The consistent chiral Fock state expansion is given by the light-cone wave function (WF) technique developed in [10] and discussed later in [11]. Here the renormalization of the axial charge by the diagrams of Fig.1a is treated as a πN Fock component contribution to the axial charge. The virtual corrections (Fig.2) do renormalize pure baryonic core, so that the whole WF is normalized to unity. The corollary of the light-cone technique is that after factoring out the angular dependence the light-cone WF, related to the πNN vertex function, depends only on one invariant variable $M_{\pi N}^2$[10]:

$$M_{\pi N}^2 = \frac{m_N^2 + k_\perp^2}{\alpha} + \frac{m_\pi^2 + k_\perp^2}{1 - \alpha}. \tag{6}$$

which is the invariant mass of the πN Fock state. In (6) α is the light-cone momentum fraction carried by the nucleon.

Introducing light-cone vectors $n_\pm = \frac{1}{\sqrt{2}}(1, 0, 0, \pm 1)$ and parametrizing the external, p, and internal, k, baryon four-momenta (Fig.1a) as $p = p_+ n_+ + p_- n_-$, $k = \alpha p_+ n_+ + \beta p_- n_- + k_\perp$, $n_\pm k_\perp = 0$ and $s = m_N^{-1}(p_+ n_+ - p_- n_-)$ after integrating over β and projecting out the A_+ component of the current yields:

$$\lambda_\pi = \frac{g_{pp\pi^0}^2}{16\pi^2}\int_0^1 \frac{d\alpha}{\alpha^2(1-\alpha)}\int_0^\infty dk_\perp^2 F_{\pi NN}^2(M_{\pi N}^2(\alpha, k_\perp^2))\frac{(1-\alpha)^2 m_N^2 - k_\perp^2}{[m_N^2 - M_{\pi N}^2(\alpha, k_\perp^2)]^2} \tag{7}$$

Evaluating the integral over α reduces eq.(7) to the dispersion integral over M^2:

$$\lambda_\pi = \frac{g_{pp\pi^0}^2}{16\pi^2}\int_{(m_N+m_\pi)^2}^\infty dM^2 F_{\pi NN}^2(M^2)$$

$$\times \frac{\frac{1}{2}(\alpha_+^2 - \alpha_-^2)(M^2 + m_N^2) - (\alpha_+ - \alpha_-)(M^2 + 3m_N^2 - m_\pi^2) + 2m_N^2\ln(\alpha_+/\alpha_-)}{(M^2 - m_N^2)^2}, \tag{8}$$

where:

$$\alpha_\pm = \left[M^2 + m_N^2 - m_\pi^2 \pm \sqrt{(M^2 + m_N^2 - m_\pi^2)^2 - 4m_N^2 M^2}\right]/2M^2.$$

The form factor $F_{\pi NN}$ cuts off large masses:

$$F_{\pi NN}^2(M_{\pi N}^2) = e^{-R_{\pi NN}^2\left[M_{\pi N}^2 - (m_N + m_\pi)^2\right]} \tag{9}$$

Direct calculation of the proton electromagnetic form factor, $F^{em}(0) = \langle p|J_+^{em}(0)|p\rangle / 2p_+$, shows that the distribution of nucleon in nucleon [9],

$$f_{N/N}(\alpha) = \frac{g^2_{\pi^0 pp}}{16\pi^2}\frac{1}{\alpha^2(1-\alpha)}\int_0^\infty dk_\perp^2 \frac{F^2_{\pi NN}(M^2_{\pi N})[m_N^2(1-\alpha)^2 + k_\perp^2]}{(m_N^2 - M^2_{\pi N})^2},$$

coincides exactly with the distribution of pion in nucleon, $f_{\pi/N}(1-\alpha)$, thus reproducing the well known Llewellyn Smith's ansatz [12]. The Llewellyn Smith's relation is crucial as it guarantees the local gauge invariance and the energy-momentum conservation. Contrary, the covariant form factor $F_{\pi NN}(u/\Lambda_N^2, t/\Lambda_\pi^2)$ which depends on both the nucleon (u) and pion (t) virtualities breaks down this relation [13]. Fine tuning of Λ_π^2 and Λ_N^2 can restore the global charge conservation, but the local gauge invariance remains badly broken. Similarly, the above covariant form factor violate the energy-momentum conservation.

The πNN coupling constant, $g_{\pi NN}$, defined at the πN system threshold, as well as the cut-off parameter, $R^2_{\pi NN}$, were originally determined in [9]. Note that different cutoffs R^2_π, R^2_K and R^2_η would be an additional source of the $SU(3)$ breaking. However, the experimental data [14] analysed with the above form factor exhibit a sort of universality: $R^2_K \approx 0.6\ (GeV/c)^{-2}$ which should be compared with $R^2_\pi \approx 0.7\ (GeV/c)^{-2}$ determined from independent measurements in [9]. We put $R^2_\eta \approx 0.6\ (GeV/c)^{-2}$. Then, for

$$\lambda_K = \frac{g^2_{p\Lambda K^+}}{16\pi^2}\int_0^1 \frac{d\alpha}{\alpha^2(1-\alpha)}\int_0^\infty dk_\perp^2 F^2_{KN\Lambda}(M^2_{K\Lambda})\frac{(m_\Lambda - \alpha m_N)^2 - k_\perp^2}{(m_N^2 - M^2_{K\Lambda})^2} \qquad (10)$$

and $\lambda_\eta = \lambda_\pi(g_{\pi NN} \to g_{\eta NN}, m_\pi \to m_\eta)$ we find numerically $\lambda_\pi = 0.008$, $\lambda_K = 0.01$ and $\lambda_\eta = 0.001$. The integral λ_η is relatively small due to the smallness of $g^2_{pp\eta} = \frac{1}{9}g^2_{p\Lambda K^+}$. The $KN\Lambda$ coupling constant determined at the $K\Lambda$ system threshold equals $g^2_{p\Lambda K^+}/4\pi \approx 6$. Notice that continuation to the kaon pole (from $M^2 = (m_\Lambda + m_K)^2$ to $M^2 = m_p^2$) gives $g^2_{p\Lambda K^+}/4\pi \approx 15$. in agreement with [15].

The WF renormalization factor is obtained from the diagram of Fig.2: $Z_N = 1 + 3\overline{\lambda}_\pi + 2\overline{\lambda}_K + \overline{\lambda}_\eta$, where $\overline{\lambda}_\pi = \int_0^1 d\alpha f_{N/N}(\alpha)$ and $\overline{\lambda}_\eta = \overline{\lambda}_\pi(g_{\pi NN} \to g_{\eta NN}, m_\pi \to m_\eta)$. To get $\overline{\lambda}_K$ from $\overline{\lambda}_\pi$ the substitution $m_N^2(1-\alpha)^2 + k_\perp^2 \to (m_\Lambda - \alpha m_N)^2 + k_\perp^2$ in $\overline{\lambda}_\pi$ should be made. Then numerically $\overline{\lambda}_\pi = 0.14$, $\overline{\lambda}_K = 0.036$, and $\overline{\lambda}_\eta = 0.005$.

It is convenient to introduce the helicity distributions $S_+(\alpha)$ defined so that $\lambda_M = \int_0^1 S_+(\alpha)d\alpha$. In the static toy-model (suppressing the isospin factor), the cancellation of the spin-flip and no-flip transitions gives $\lambda_\pi = -1/3$, which can easily be deduced from the $\frac{1}{2} \otimes 1 = \frac{1}{2}$ Clebsch-Gordan coefficients. In the relativistic case the spin-flip is due to $k_\perp \neq 0$, eq.(7), and the cancellation pattern is sensitive to the form factor radius. Namely, spin-flip dominates at large α where $S_+(\alpha) < 0$ (Fig.3). However, at smaller α, when significant fraction of the nucleon's momentum is carried by the pion, one finds $S_+(\alpha) > 0$. Notice, that this is the relativistic recoil region which is not treated properly in the non-relativistic calculations.

The spin-flip and non-flip contributions to $S_+(\alpha)$ are cancelled in λ_π resulting in $\lambda_\pi \approx 0.008$ at $R^2 = 0.7\ (GeV/c)^{-2}$. The integral $\lambda_\pi(R^2)$ depends strongly on the cutoff parameter R^2 (Fig.4). At $R^2 \approx 0.4\,(GeV/c)^{-2}$, corresponding to the momentum cutoff $k_{\perp c} \approx R^{-1}[m_\pi/(m_N + m_\pi)]^{1/2} \approx 0.6\,GeV/c$, $\lambda_\pi(R^2)$ becomes negative and steeply falls at smaller R^2 (bigger $k_{\perp c}$). These results contradict to non-relativistic intuition. In the Cloudy Bag Model λ_π stays negative valued independently of the bag radius.

Notice that the quantity $\lambda_K(R^2)$ closely related to the polarized strange quark density,

$$\Delta s = (2/9)(5F + D)\lambda_K Z_N^{-1}, \qquad (11)$$

remains almost unchanged (Fig.4). To derive eq.(11) it suffices to make use of the $SU(3)$ symmetry of the tree level baryon wave functions and to consider the relevant $K-$exchange diagrams. Throughout this paper the $KN\Lambda$ and $KN\Sigma$ couplings as well as pion-baryon ones are supposed to be related via SU(6).

The axial charge $g_A^8(R^2)$ decreases rapidly with decreasing R^2: in the exact $SU(6)$ limit $g_A^8(R^2) \approx [1 + 3\lambda_\pi(R^2)]Z_N^{-1}$. The renormalized flavour-singlet axial charge, $g_A^0 = g_A^8 + 3\Delta s$, can reach rather small values at $k_{\perp c} > 1\,GeV/c$ ($g_A^0 \approx 0.25$ at $k_{\perp c} \approx 1.5\,GeV/c$). However, at large $k_{\perp c}$ the renormalized $\beta-$decay coupling $g_A^3 \approx \frac{5}{3}[1 - \lambda_\pi(R^2)]Z_N^{-1}$ becomes uncomfortably small ($g_A^3 \approx 0.5$) due to the growing WF renormalization factor, Z_N. Thus the resolution of the spin crisis problem proposed in [16] and basing on the allowance for the large pion momentum cutoffs looks doubtful. Note in this connection that the intermediate spin-$\frac{3}{2}$ decuplet contributions (see below) to both g_A^8 and g_A^3 are large and positive.

3. In sect.2 we made use of the Sudakov (light-cone) variables to reduce the one-loop integrals. It can easily be checked that the non-covariant time-ordered perturbation theory in the Infinite Momentum Frame (IMF) leads to exactly the same results for both the electromagneic form factors and renormalized axail couplings. In IMF one operates with on-mass shell particles (off-energy shell intermediate states). This simplifies calculations with spin-$\frac{3}{2}$ decuplet significantly and allows to get rid of ambiguities with the off-mass shell behaviour of spin-$\frac{3}{2}$ particles (spin-1/2 admixture) [17].

The momentum space propagator of the decuplet field $D_\mu(x)$ can be taken in the form:

$$\Delta_{\lambda\rho}(k) = \frac{\hat{k} + m_\Delta}{k^2 - m_\Delta^2 + i\epsilon}\left(\frac{2k_\lambda k_\rho}{3m_\Delta^2} - g_{\lambda\rho} + \frac{1}{3}\gamma_\lambda\gamma_\rho - \frac{k_\lambda\gamma_\rho - k_\rho\gamma_\lambda}{3m_\Delta}\right).$$

The $SU(3)$ indices are suppressed. The decuplet axial current A^a_μ is written down as $A^a_\mu = CT^a\bar{D}_\nu\gamma_\mu\gamma_5 D_\nu$. For the on-shell delta the term $\frac{1}{2}\epsilon_{\mu\nu\lambda\rho}\bar{D}_\nu\gamma_\lambda D_\rho$, which could contribute to the off-shell delta current A^a_μ, can be safely neglected.

Normalization is such that: $2\langle\Delta^{++}|A^3_\mu|\Delta^{++}\rangle = 2Cm_\Delta s_\mu$. The octet ($B$) -decuplet ($D$) interaction with the pseudo-scalar meson field $M(x)$ is described by $\mathcal{L}_{BDM}(x) = ig_{BDM}\bar{B}(x)D_\mu(x)\partial_\mu M(x) + h.c.$ The one-loop correction to g_A^8 caused by the intermediate $\pi\Delta$ states (Fig.1b) is $2C\rho_\pi/Z_N$. The integral flow of the proton helicity, ρ_π, for the $\pi^-\Delta^{++}$ Fock component is as follows:

$$\rho_\pi = \frac{g^2_{\pi^+ p\Delta^{++}}}{16\pi^2}\int_0^1 \frac{d\alpha}{\alpha^2(1-\alpha)}\int_0^\infty dk_\perp^2 F^2_{\pi N\Delta}(M^2)\frac{a_+^{DD}}{(m_N^2 - M^2)^2}, \qquad (12)$$

where

$$a_+^{DD} = \frac{1}{18\alpha^2 m_\Delta^2}\Big\{(m_\Delta + \alpha m_N)^2(m_\Delta^2 - \alpha^2 m_N^2)^2 + k_\perp^2(m_\Delta + \alpha m_N)^2$$
$$\times\Big[(m_\Delta + \alpha m_N)^2 + 4\alpha m_N m_\Delta\Big] - (k_\perp^2)^2\Big[(m_\Delta - \alpha m_N)^2 - 4\alpha m_N m_\Delta\Big] - (k_\perp^2)^3\Big\} \qquad (13)$$

Here α is the IMF couterpart of the Sudakov variable introduced above: in IMF $p = (P + m_N^2/2P, \vec{0}, P)$, $k = (\alpha P + (m_\Delta^2 + k_\perp^2)/2\alpha P, \vec{k}_\perp, \alpha P)$ and $s = m_N^{-1}(P, \vec{0}, P + m_N^2/2P)$

at $P \to \infty$. For $R^2_{\pi N\Delta} = 1.4\,(GeV/c)^{-2}$ determined in[9] $\rho_\pi = 0.042$. Intermediate $\pi\Delta$ states contribute positively to the proton helicity distribution (Fig.3) thus increasing the axial isosinglet coupling, $g^0_A = [0.60(3F) - 0.61D + 0.084\,C]/Z_N$.

The axial current $A^a_\mu = H\,T^a\,\bar{B}D_\mu$ mediating transitions between octet and decuplet is normalized as follows $2\left\langle p|A^3_\mu|\Delta^+\right\rangle = 2H\sqrt{m_\Delta m_N}\cdot s_\mu$. Then the interference of the πN and $\pi\Delta$ Fock states (Fig.1c) results in

$$\nu_\pi = \frac{g_{pp\pi^0}g_{\pi^+p\Delta^{++}}}{16\pi^2}\int_0^1 \frac{d\alpha}{\alpha^2(1-\alpha)}\int_0^\infty dk_\perp^2 F_{\pi NN}(M^2_{\pi N})F_{\pi N\Delta}(M^2_{\pi\Delta})$$
$$\times\frac{(m^{*2}-\alpha^2 m_N^2)^2 - k_\perp^2(k_\perp^2 - 2\alpha m_N m^*)}{3\alpha m^*\,[m_N^2 - M^2_{\pi N}]\,[m_N^2 - M^2_{\pi\Delta}]} \qquad (14)$$

For consistency we put (within the loop) $m_\Delta = m_N = m^*$ and $m^* = 1\,(GeV/c^2)$. Numerically $\nu_\pi = 0.05$. Calculation of the $\pi\Delta$ contribution to the proton electromagnetic form factor, shows again that the form factor $F(M^2)$ ensures the local charge conservation and the energy-momentum sum rules. The distribution of delta in nucleon [9],

$$f_{\Delta/N}(\alpha) = \frac{g^2_{\pi^+p\Delta^{++}}}{96\pi^2 m_\Delta^2}\frac{1}{\alpha^4(1-\alpha)}$$
$$\times\int_0^\infty dk_\perp^2 \frac{F^2_{\pi N\Delta}(M^2_{\pi\Delta})[(m_\Delta+\alpha m_N)^2 + k_\perp^2]^2[(m_\Delta - \alpha m_N)^2 + k_\perp^2]}{[m_N^2 - M^2_{\pi\Delta}]^2} \qquad (15)$$

coincides with the distribution of pion in nucleon, $f_{\pi/N}(1-\alpha)$.

The integral flux of $\pi\Delta$ states entering the WF renormalization factor Z_N equals $2\bar{\rho}_\pi$. From Fig.2b it follows that in terms of $f_{\Delta/N}$: $\bar{\rho}_\pi = \int_0^1 d\alpha f_{\Delta/N}(\alpha)$.

For the numerical estimates the tree level axial couplings are supposed to be related via $SU(6)$, $F/D = 2/3$. We put $D = 1$ and $C = 3$, $H = 4\sqrt{2}/3$ as it follows from the nonrelativistic quark model [18]. Then the renormalized polarized quark densities are $\Delta u = 0.98$, $\Delta d = -0.24$, $\Delta s = 0.006$ and the renormalized flavour -singlet axial coupling is $g^0_A = 0.75$.

4. As a consistency check we calculate here the ratios of the axial-vector and vector form factors measured in semileptonic decays of hyperons. These ratios parametrized in the Cabibbo framework were found to be [19]:

$$(g_A/g_V)_{\Xi^-\to\Lambda} = F - D/3 = +0.25 \pm 0.05,$$
$$(g_A/g_V)_{\Sigma^-\to n} = F - D = -0.34 \pm 0.05,$$
$$(g_A/g_V)_{\Lambda\to p} = F + D/3 = +0.70 \pm 0.03.$$

From the measured value of the branching ratio $R = \Gamma(\Sigma^- \to \Lambda e^-\bar{\nu})/\Gamma(\Sigma^- \to all)$ it was found [19]: $D = 0.720 \pm 0.020$. The value of $(g_A/g_V)_{n\to p}$ [20] is: $F + D = 1.2573 \pm 0.0028$. Radiative corrections modify the tree level relations significantly:

$$(g_A)_{n\to p} = 0.58\,F + 0.58\,D + 0.027\,C + 0.095\,H = 1.22,$$
$$(g_A)_{\Xi^-\to\Lambda} = 0.74\,F - 0.75\,(D/3) + 0.01\,C - 0.004\,H = 0.26,$$
$$(g_A)_{\Lambda\to p} = 0.64\,F + 0.61\,(D/3) + 0.018\,C + 0.05\,H = 0.77,$$

$$(g_A)_{\Xi^-\to\Lambda} = 0.71\,D - 0.02\,F + 0.01\,C + 0.02\,H = 0.76,$$
$$(g_A)_{\Sigma^-\to n} = 0.68\,F - 0.66\,D + 0.006\,C - 0.035\,H = -0.26$$

Here we refer to g_A instead of g_A/g_V. The agreement seems to be reasonable, especially if we take into account that the assumption of the tree level $SU(6)$ symmetry left us without any adjusting parameters.

5. Naively, one could have expected rather large flow of nucleon's spin to the angular momentum of the pionic cloud. Our analysis shows that this is not a case. One of the major reasons is that the πN Fock component of the nucleon receives large contribution from beyond the conventional static limit. Besides the $\pi\Delta$ Fock states contribute significantly and Δ tends to retain proton's polarization. As a result the renormalization of the flavour-singlet coupling g^0_A by the radiative mesonic corrections is very weak: $(g^0_A)_{SU(6)} = 1$, we have found $g^0_A = 0.75$. This implies that the axial anomaly does indeed play a crucial role in understanding the EMC spin crisis [3].

We found that the assumption of exact $SU(6)$ symmetry for bare axial couplings is a good starting point for quantitative analysis. Renormalized axial couplings are in good agreement with measured rates of the Gamov-Teller transitions in the hyperon octet.

We notice that the one-loop corrections, originate rather specific spin structure of the proton. The light-cone proton helicity distribution is oscillating function of the momentum fraction α carried by intermediate nucleon. Spin-flip and non-flip contributions are cancelled in the integral for the proton axial charge. However the integral flux of pions is not small, $3\bar{\lambda}_\pi \approx 0.4$, and renormalization of axial charges is mainly due to the WF renormalizaton. Intermediate Δ contributes positively to the proton helicity and significantly increases the flux of pions $2\bar{\rho}_\pi \approx 0.2$.

The polarized strange quark density is compatible with zero $\Delta s \approx 0.006$ and depends on the cutoff parameter rather weakly. The integral flux of nonperturbative $s\bar{s}$ pairs is estimated as $2\bar{\lambda}_K \approx 0.07$. In our approach the contribution of intermediate Σ^* to Δs as well as the interference term $\Sigma^* - \Sigma$ are suppressed. We shall expand on this subject elsewhere.

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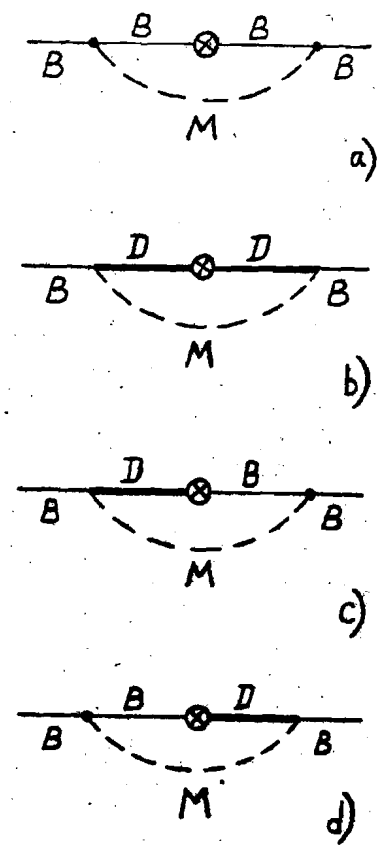

Fig.1 The diagrams contributing to the renormalization of the baryon octet axial currents. The circled cross represents an insertion of the axial current A_μ. The π, K, η–fields are marked by the symbol M. The decuplet ones are marked by the symbol D.

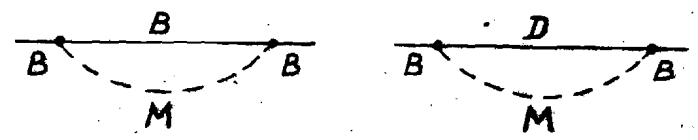

Fig.2 Wavefunction renormalization diagrams.

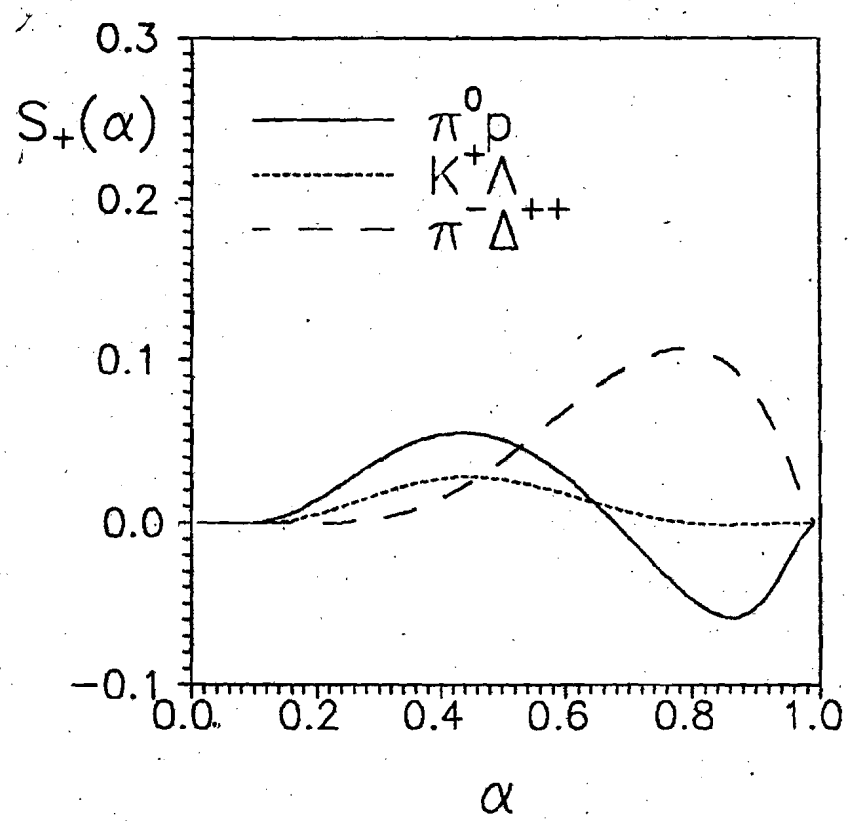

Fig.3 The proton helicity (" + " -component of the axial charge) against the Sudakov variable α (the proton momentum fraction carried by intermediate baryon). Contributions of the $\pi^0 p$, $K^+\Lambda$ and $\pi^-\Delta^{++}$ Fock components of the proton are shown.

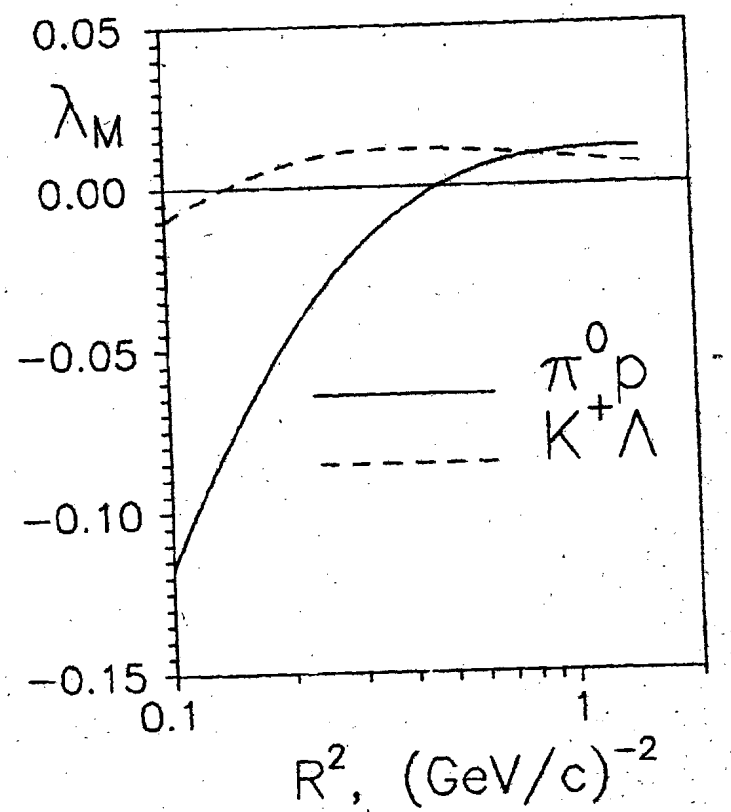

Fig.4 R^2 -dependence of the integral helicity flow $\lambda_M(R^2)$ for both the strange, $K^+\Lambda$, and non-srange, $\pi^0 p$, Fock components of the ptoton.

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