Transverse $Λ^0$ polarization in inclusive quasi-real photoproduction at the current fragmentation

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Abstract. It is shown that the recent HERMES data on the transverse $Λ^0$ polarization in the inclusive quasi-real photoproduction at $x_F > 0$ can be accommodated by the strange quark scattering model. Relations with the quark recombination approach are discussed.

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1 Introduction

Polarization of $Λ^0$ hyperons has been under scrutiny almost since the very moment of its discovery. Investigations of the phenomenon received especially great impetus in 1976 due to the striking experimental results obtained at FERMILAB, where the hyperons produced in $p$N collisions at 300 GeV proton beam energy were highly polarized [1]. The polarization was transverse and negative, directed opposite to the unit vector $\mathbf{n} \propto [\mathbf{p}_b \times \mathbf{p}_A]$, where $\mathbf{p}_b$ and $\mathbf{p}_A$ are the beam and hyperon momenta, respectively.

Only this direction is allowed by the parity conservation in strong interactions provided the incident particle’s spin-flavor part of the wave function to be combined of the $ud$ diquark in a singlet spin state and the strange quark of spin $1/2$, or formally $|A\rangle_{1/2} = |ud\rangle_0 |s\rangle_{1/2}$, where the subscriptions denote the spin states. Therefore, the total spin of $Λ^0$ is entirely given by the spin of its valence $s$ quark. There is also an alternative look at the spin transfer in fragmentation appeared after publication of the polarized deep inelastic lepton-nucleon scattering (DIS) data of the EMC Collaboration [17]. It suggests that the spin carried by the valence quarks is only a part of the total nucleon spin, the rest one being attributed, for example, to the orbital angular momenta of the valence quarks and to the nucleon sea (sea quarks, antiquarks and gluons). Which picture, SU(6) or DIS, is suitable for the description of the process remains still an important issue [18, 19, 20, 21, 22, 23]. The $Λ^0$ hyperon can provide here a useful instrument for the study of the spin effects in strong interactions.

We used the SU(6) approach throughout this paper. The choice was dictated by a wish to keep the quark scattering model as it is, encouraged by the SU(6) based calculations of the longitudinal $Λ^0$ polarization in $e^+e^−$ annihilation at the $Z^0$ pole [24] and their successful experimental verification [25, 26].

According to the empirical rules proposed by DeGrand and Miettinen, the polarization sign depends on whether the $s$ quark is accelerated (increases its energy) or decelerated (decreases its energy) in the $Λ^0$ formation process [5]. To illustrate, there are no valence $s$ quarks in the initial state of the $pp$ reaction so that they come from the quark sea to form the final $Λ^0$. But the sea quarks predominantly populate small $x$-states ($x$ is Bjorken variable) and consequently increase their average energy coming in the valence content of $Λ^0$. Here the polarization is negative. Contrary, incident pseudoscalar kaons of the $K^−p$ reaction already contain valence strange quarks mostly decelerated in the hadronization because the average energy of the created $Λ^0$-s is less than that of the $K^−$ beam. In this
case the sign is positive. Similar ideas were implemented in flux-tube models with orbital angular momentum [4]. It was natural to wonder whether the polarization took place in $A^0$ electro and/or photoproduction. Would one observe here the same features with hadron-hadron reactions? These questions have been investigated, for example, in experiments on high energy $\gamma N$ scattering performed at CERN [27] and SLAC [28] in the beginning of the eighties. However, their statistical accuracy is indecisive and would hardly enable one to conclude on the magnitude or on the sign of the polarization.

The transverse $A^0$ polarization has been also measured in unpolarized $e^+e^-$ annihilations, for example, by TASSO Collaboration at 14, 22 and 34 GeV center-of-mass (cms) energies [29] and near the $Z^0$ pole by OPAL [26]. The polarization observed in both experiments is consistent with zero. Practically, this process can be a good place for deriving some important information on the hadronization phase. In particular, it can assist understanding in which extent final state interactions contribute to the transverse polarization [30].

In light of the scarce statistics for the $A^0$ photoproduction process, the HERMES experiments on the 27.6 GeV positron beam scattering off the nucleon target acquires a particular status providing a good opportunity for observation of the polarization in electroproduction. The collaboration has measured nonzero positive transverse polarization, when most of the intermediate photons are quite close to the mass shell, i.e. $Q^2 = -(p_{ei} - p_{ef})^2 \approx 0$ GeV$^2$ [31], where $p_{ei}$ and $p_{ef}$ are the 4-momenta of the initial and scattered electrons, respectively (quasi-real photoproduction).

Experimental properties of the polarization at HERMES turned out to be very reminiscent of those in the $K^-p$ reaction [4], which has been successfully described by a model assuming the polarization to appear mostly via strange quark scattering in a color field [6, 7]. Thus, there are indications that mechanisms responsible for the phenomenon in the $K^-p$ reaction, while $ep$ may be similar, at least, within the covered kinematic region. These arguments inspired us to apply the model to the $A^0$ quasi-real photoproduction data obtained by HERMES.

Another goal of this paper is qualitatively to discuss some relations between the calculations presented herein and the quark recombination model (QRM) [14].

2 Quark scattering model for $A^0$

Electrons, when scattering off nuclei, have been known to be able to acquire polarization. Analytically, it can be found within QED by considering a process of Dirac pointlike particle scattering off static Coulomb potential belonging to next-to-leading order amplitudes are taken into account [32,33] (see also Ref. [34]). The corresponding formula reads

$$P = \frac{2\alpha_{em}mp}{E^2} \left[ \sin^2 \theta / 2 \ln(\sin \theta / 2) \right] \cos \theta / 2 n, \quad (1)$$

where $E$, $p$, $m$ and $\theta$ are the energy, momentum by magnitude, mass and scattering angle of the electron, respectively, $\alpha_{em}$ is the fine structure constant, $n = [p_i \times p_f]$, $p_i$ and $p_f$ are the vectors of the electron momenta in the initial ($i$) and final ($f$) states.

In Ref. [6], Szwed proposed to explain the $A^0$ polarization as polarization of its valence strange quark in scattering using Eq. (1). The $A^0$ wave function favors such a consideration quite well. The idea is to perform the following interchanges in Eq. (1): electron ↔ quark, $\alpha_{em} \rightarrow C\alpha_s$ (Coulomb potential ↔ color field), where $\alpha_s$ is the strong coupling and $C$ is the color factor.

This approach has been applied to describe the polarization in the $K^-p$ reaction and successfully reproduced its main features at $2C\alpha_s = 5.0$ and the $s$ quark mass $m_s = 0.5$ GeV [7].

Due to some peculiarities of the HERMES experiment, the polarization was not measured in the traditional form of the dependence on $x_F = 2p_z/\sqrt{s}$ ($p_z$ is the longitudinal component of the detected particle momentum, $\sqrt{s}$ is the total cms energy). Therefore, we have expressed the model in terms of the light cone variable $\zeta$ the available data depend on [31]. It is defined as

$$\zeta(f) = \frac{E(f) + p_z(f)}{E_b + p_z(b)}, \quad (2)$$

here the subscription $b$ denotes the beam. We note that $\zeta$ is invariant under Lorentz boosts being useful in its application.

According to recipes given in Ref. [7], one should move to a frame, where the magnitudes of the initial and final $s$ quark momenta are the same (originally called S-frame). It is reached by performing a Lorentz transformation along the proton momentum. For this purpose, one can write

$$(p_i \cdot p_f) = p^2 (1 - \cos \theta) + m^2_s, \quad (3)$$

$$p_T f = p_T r = p \sin \theta, \quad (4)$$

where $p_{i(f)}$ are the 4-momenta of the scattering quark, $p_T f$ is the transverse momentum of the scattered quark in the center-of-mass frame of the $K^-p$ reaction, while $p = \sqrt{E^2 - m^2_s}$, $p_T$ and $\theta$ refer to the S-frame.

On the other hand, using Eq. (2) leads to

$$(p_i \cdot p_f) - m^2_s = \frac{m^2_s}{2} \frac{(\zeta_i - \zeta_f)^2}{\zeta_i \zeta_f}$$

$$+ \frac{1}{2} \left[ p_{T i} \frac{\zeta_f}{\zeta_i} + p_{T f} \frac{\zeta_i}{\zeta_f} \right] + (p_{T i} \cdot p_{T f}), \quad (5)$$

where $(p_{T i} \cdot p_{T f})$ is the ordinary scalar product of the transverse momentum vectors.

Neglecting, as a first approximation, the incident quark transverse momentum ($p_{T i} = 0$), after some algebra, one can obtain from Eqs. (5) that

$$\cos \theta / 2 = \frac{\xi V^2}{(1 - \xi^2) + V^2}, \quad (6)$$

where $V = (\zeta_i - \zeta_f)/\zeta_i \zeta_f$, $\xi = \sqrt{2E/E_b}$, $E$ is the electron energy. If $E = E_b$, we get $\cos \theta / 2 = (1 - \xi^2)/2$. This formula agrees well with measured data [31].
\[ V = \frac{(1 - \xi)^2 + V^2}{2 \sqrt{\xi} \sqrt{(1 - \xi)^2 + (1 - \xi)V^2}}. \] (7)

with the variables \( V(T) \), and \( \xi \) defined by

\[ V(T) = \frac{p(T)}{m_s}, \quad \xi = \frac{\zeta_f}{\zeta_i}. \] (8)

By using relations (6) and (7), Eq. (1) can be rewritten as

\[ P(\xi, V_T) = -\frac{2C_{\alpha_s} V}{1 + V^2} \frac{\sin^2 \theta/2 \ln(\sin \theta/2)}{\cos \theta/2}. \] (9)

Note that the minus sign in Eq. (9) appeared to satisfy the rule of DeGrand and Miettinen in the region of our interest (\( \xi < 1 \)).

3 Calculations and results

Our calculations concern the polarization in the region of \( 0.25 \lesssim \zeta \lesssim 0.5 \), which, according to Ref. [31], corresponds to the events of \( x_F > 0 \) in the cms frame of the \( \gamma^* p \) reaction (current fragmentation). The procedure is now straightforward. In the model discussed above, one substitutes the \( K^- \) meson by the intermediate quasi-real photon \( \gamma^* \). Since the considered hyperons are produced in the photon fragmentation region, we assumed that the \( A^0 \) kinematic is determined here in the main by the quarks originated from the photon. The \( (ud)_0 \) diquarks are supposed to come mostly from the proton target. The polarization process is schematically shown in Fig. 1. Note that a similar diagram, corresponding to the \( s \) quark scattering off the scalar diquark \( (ud)_0 \), appears also in the QRM [14] when one calculates the polarization in the \( K^- p \) reaction, however the interaction is chosen to be scalar.

![Schematic diagram of the \( A^0 \) polarization process](image1)

**Fig. 1.** A schematic diagram of the \( A^0 \) polarization process in quasi-real photoproduction. The \( s \) quarks originated from the photon scatter off the target color field, getting thus polarized, and form the final hyperon recombining with a \( (ud)_0 \) diquarks from the proton target. An arrow over a letter indicates the polarization.

Within the present approach, one needs to know the \( \zeta_i \) distribution of the incident \( s \) quarks originated from the quasi-real photons emitted by the positron beam. To find it, we used the PYTHIA 6.2 program [35], adopting thus the positron-to-quark transition mechanisms implemented therein. The obtained distribution is shown in Fig. 2 (scattered plot) together with its fit (solid line).

![Graphical representation of \( \zeta_i \) distribution](image2)

**Fig. 2.** The \( \zeta_i \) distribution of the quarks originated from the positron beam according to the PYTHIA program (scattered plot). The corresponding fit is presented by the solid line. Events assumed to contribute to the region of \( \zeta > 0.25 \) are hatched.

The final \( s \) quark kinematic was determined according to the following

\[ \zeta_f = \frac{m_s}{m_A} \zeta_A, \quad V_T = \frac{p_{TA}}{m_A}, \] (10)

where \( \zeta_A \) and \( p_{TA} \) refer to the detected \( A^0 \) hyperons. Let us omit in the sequel the index \( A \). The relations in Eq. (10) define the quark momenta as fixed fractions of the corresponding final hyperon momenta. In particular, the direction of the \( s \) quark transverse momentum is assumed to point in most cases in the direction of \( p_{TA} \) [6]. It is, of course, only an approximate picture. One should take into account that the quarks are just constituents of the hyperon but not free. In fact, a momentum component of a quark inside a hadron is not fixed but has an intrinsic distribution.

Note that there should be a threshold for the \( A^0 \) production in the region of \( \zeta > 0.25 \) initiated by the incoming quarks due to the undetected hadron system always exists and carries also away some part of the energy. In other words, not all the \( \zeta_i \) events will contribute to the \( \zeta > 0.25 \) region. Thus, in the calculations, we took into account only those quarks with \( 0.5 \leq \zeta_i \leq 1 \) (hatched area in Fig. 2).

As the experimental \( \zeta \) dependence of the polarization is available integrally over \( p_T \), we additionally introduced averaging over the transverse momentum of the detected \( A^0 \)-s. In this case, the polarization was determined as

\[ P_\zeta = \int d\zeta_i dp_T \ h(p_T) \ P \left( \frac{\zeta}{\zeta_i}, p_T \right) f(\zeta_i). \] (11)

Here, \( h(p_T) \) and \( f(\zeta_i) \) are the \( p_T \) and \( \zeta_i \) distribution functions of the detected \( A^0 \)-s and the incident quarks.
respective, \( P(\zeta/\zeta_\alpha, p_T) \) is the polarization defined by Eq. (9).

For similar reasons, we determined the \( p_T \) dependence of the polarization as

\[
P_{p_T} = \int d\zeta d\zeta' \ g(\zeta) P \left( \frac{\zeta}{\zeta'}, p_T \right) f(\zeta),
\]

where \( g(\zeta) \) is the \( \zeta \) distribution function of the detected hyperons.

Using Eqs. (11) and (12), we carried out the calculations. We used a typical \( \zeta \) distribution of the detected \( \Lambda^0 \)s \((0.25 \leq \zeta \leq 0.5)\). For \( h(p_T) \), we adopted that obtained by HERMES \((0.2 \text{ GeV} \leq p_T \leq 1.2 \text{ GeV})\) [36]. Note that all the distributions were prenormalized to unity. As the free parameter values, we have taken \( 2C\alpha_s = 5.0 \) and \( m_s = 0.5 \) GeV.

The numerical results in comparison with the HERMES data are shown in Fig. 3. One can see that the experiment is reasonably reproduced.

\[\text{Fig. 3. Numerical results (lines) in comparison with the HERMES data (solid points). The } \zeta(\zeta_T) \text{ dependence of the } \Lambda^0 \text{ polarization is presented in the upper (lower) panel. The } p_T \text{ data are taken from Ref. [31].}\]

4 Conclusion

The results obtained herein should be regarded only as qualitative. We neglected the transverse momentum of the incident quarks \((p_T)\), while, certainly, a strict consideration would require taking it into account. However, for the goal, the present work aimed at, such an approximation reflects the general tendencies. It would be fair to expect a relatively narrow \( p_T \) distribution for the events contributing to the region of \( \zeta > 0.25 \). To find the \( \zeta \) distribution, we used the PYTHIA program, which gives, in turn, rather qualitative than quantitative predictions. We did not take the contributions from the heavier resonances, such as \( \Sigma^0 \), \( \Xi \) and \( \Sigma^+ \), into account, their values are presumably considerable in \( \Lambda^0 \) polarization [24, 37, 38]. A difficulty is also caused by the impossibility to derive the running coupling constant \( \alpha_s \) from the HERMES data.

In fact, the calculations by Eqs. (11) and (12) have been carried out similarly with the quark recombination model [14], the latter describes the polarization in a more quantitative manner. In the QRM, the central point is the squared subprocess amplitude averaged over the Bjorken variables in the initial and final states, their role, in our case, were played by \( \zeta \) and \( \zeta \). A substantial distinction between the QRM and the present quark scattering approach (QSM) is the interaction. For the QRM, it has been assumed to be a scalar force, while the QSM calculations are based on QCD. Thus it seems to be attractive alternatively to specify the QRM interaction by the color one. Doing it could be regarded as a further development of the present approach. It will include, in particular, the transverse momentum of the incident quarks, the structure functions of the projectile as well as the outgoing hyperon instead of the approximations of Eq. (10), it will also more explicate the underlying subprocesses. An estimation of the contributions of the heavier resonances is in progress now.

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List of Figures

1 A schematic diagram of the $Λ^0$ polarization process in quasi-real photoproduction. The $s$ quarks originated from the photon scatters off the target color field, getting thus polarized, and form the final hyperon recombining with a $(ud)_0$ diquarks from the proton target. An arrow over a letter indicates the polarization. . . . . . . . . .

2 The $ζ$ distribution of the quarks originated from the positron beam according to the PYTHIA program (scattered plot). The corresponding fit is presented by the solid line. Events assumed to contribute to the region of $ζ > 0.25$ are hatched. . . . . . . . . .

3 Numerical results (lines) in comparison with the HERMES data (solid points). The $ζ (p_T)$ dependence of the $Λ^0$ polarization is presented in the upper (lower) panel. The data are taken from Ref. [31]. . . . . . . . .