Chiral quark filtering mechanism of hyperon polarization

S. M. Troshin, N. E. Tyurin

Institute for High Energy Physics,
Protvino, Moscow Region, 142281, Russia

Abstract

The model combined with unitarity and impact parameter picture provides a rather simple mechanism for generation of hyperon polarization in collision of unpolarized hadrons. We concentrate on a particular problem of Λ-hyperon polarization and derive its linear $x_F$-dependence as well as its energy and transverse momentum independence at large $p_\perp$ values. Mechanism responsible for the single–spin asymmetries in pion production is also discussed.
Introduction

One of the most interesting and persistent for a long time spin phenomena was observed in inclusive hyperon production in collisions of unpolarized hadron beams. A very significant polarization of $\Lambda$–hyperons has been discovered almost three decades ago [1]. There is a list of theoretical models which relate polarization mechanism with various aspects of hadron interaction dynamics [2] but till now it has not obtained a satisfactory explanation. Experimentally the process of $\Lambda$-production has been studied more extensively than other hyperon production processes. Therefore we will emphasize on the particular riddle of $\Lambda$–polarization because spin structure of this particle is most simple and is determined by strange quark only. We also provide comments on how this mechanism can be used for the explanation of single-spin asymmetries in the inclusive pion production.

It should be noted that understanding of transverse single-spin asymmetries in DIS (in contrast to the hyperon polarization) has observed significant progress during last years; this progress is related to an account of final-state interactions from gluon exchange [3, 4] – coherent effect not suppressed in the Bjorken limit.

Experimental situation with hyperon polarization is widely known and stable for a long time. Polarization of $\Lambda$ produced in the unpolarized inclusive $pp$–interactions is negative and energy independent. It increases linearly with $x_F$ at large transverse momenta ($p_\perp \geq 1$ GeV/c), and for such values of transverse momenta is almost $p_\perp$-independent [1].

On the theoretical side, perturbative QCD with a straightforward collinear factorization scheme leads to small values of $\Lambda$–polarization [5, 6] which are far below of the corresponding experimental data. Modifications of this scheme and account for higher twists contributions allows to obtain higher magnitudes of polarization but do not change a decreasing dependence proportional to $1/p_\perp$ [7, 8, 9] at large transverse momenta. It is difficult to reconcile this behavior with the flat experimental data dependence on the transverse momenta. Inclusion of the internal transverse momentum of partons ($k_\perp$–effects) into the so called polarizing fragmentation functions leads to similar decreasing polarization [10]. In addition it should be noted that the perturbative QCD has also problems in the description of the unpolarized scattering, e.g. in inclusive cross-section for $\pi^0$-production, at the energies lower than the RHIC energies [11].

The essential point of the approaches mentioned above is that the vacuum at short distances is taken to be a perturbative one. There is another possibility. It might happen that the hyperon polarization dynamics originates from the genuine nonperturbative sector of QCD (cf. e.g. [12]). The point of view that the polarization has its roots hidden in the nonperturbative sector of QCD is not an isolated one and several approaches based on nonperturbative dynamics has appeared up to now. We briefly mention them later.
In the nonperturbative sector of QCD the two important phenomena, confinement and spontaneous breaking of chiral symmetry ($\chi$SB)\(^{[13]}\) should be reproduced. The relevant scales are characterized by the parameters $\Lambda_{QCD}$ and $\Lambda_{\chi}$, respectively. Chiral $SU(3)_L \times SU(3)_R$ symmetry is spontaneously broken at the distances in the range between these two scales. The $\chi$SB mechanism leads to generation of quark masses and appearance of quark condensates. It describes transition of current into constituent quarks. Constituent quarks are the quasiparticles, i.e. they are a coherent superposition of bare quarks, their masses have a magnitude comparable to a hadron mass scale. Therefore hadron is often represented as a loosely bounded system of the constituent quarks. These observations on the hadron structure lead to understanding of several regularities observed in hadron interactions at large distances. It is well known that such picture provides reasonable values for the static characteristics of hadrons, for instance, their magnetic moments. The other well known direct result is appearance of the Goldstone bosons.

The most recent approach to single–spin asymmetries (SSA) based on nonperturbative QCD has been developed in \(^{[14]}\) where, in particular, $\Lambda$-polarization has been related to the large magnitude of the transverse flavor dipole moment of the transversely polarized baryons in the infinite momentum frame. It is based on the parton picture in the impact parameter space and assumed specific helicity–flip generalized parton distribution.

The instanton–induced mechanism of SSA generation was considered in \(^{[15,16]}\) and relates those asymmetries with a genuine nonperturbative QCD interaction. It should be noted that the physics of instantons (cf. e.g. \(^{[17]}\)) can provide microscopic explanation for the $\chi$SB mechanism\(^{1}\).

We discuss here mechanism for hyperon polarization based on chiral quark model\(^{2}\)\(^{[13]}\) and the filtering spin states related to unitarity in the $s$-channel. This mechanism connects polarization with asymmetry in the position (impact parameter) space.

1 Chiral quark model and the mechanism of spin states filtering

As it was already mentioned constituent quarks and Goldstone bosons are the effective degrees of freedom in the chiral quark model. We consider a hadron consisting of the valence constituent quarks located in the central core which is

\(^{1}\)We are grateful to Dmitri Diakonov for the interesting communication on this matter regarding the polarization phenomena.

\(^{2}\)It has been successfully applied for the explanation of the nucleon spin structure \(^{[13]}\).
embedded into a quark condensate. Collective excitations of the condensate are the Goldstone bosons and the constituent quarks interact via exchange of Goldstone bosons; this interaction is mainly due to a pion field which is of the flavor– and spin–exchange nature. Thus, quarks generate a strong field which binds them.

At the first stage of hadron interaction common effective self-consistent field is appeared. Valence constituent quarks are scattered simultaneously (due to strong coupling with Goldstone bosons) and in a quasi-independent way by this effective strong field. Such ideas were already used in the model which has been applied to description of elastic scattering and hadron production.

The initial state particles (protons) are unpolarized. It means that states with spin up and spin down have equal probabilities. The main idea of the proposed mechanism is the filtering of the two initial spin states of equal probability due to different strength of interactions. The particular mechanism of such filtering can be developed on the basis of chiral quark model, formulas for inclusive cross section (with account for the unitarity) and notion on the quasi-independent nature of valence quark scattering in the effective field.

We will exploit the feature of chiral quark model that constituent quark $Q_\uparrow$ with transverse spin in up-direction can fluctuate into Goldstone boson and another constituent quark $Q'_\downarrow$ with opposite spin direction, i.e. perform a spin-flip transition:

$$Q_\uparrow \rightarrow GB + Q'_\downarrow \rightarrow Q + \bar{Q'} + Q'_\downarrow. \tag{1}$$

An absence of arrows means that the corresponding quark is unpolarized. To compensate quark spin flip $\delta S$ an orbital angular momentum $\delta L = -\delta S$ should be generated in final state of reaction. The presence of this orbital momentum $\delta L$ in its turn means shift in the impact parameter value of the final quark $Q'_\downarrow$ (which is transmitted to the shift in the impact parameter of $\Lambda$)

$$\delta S \Rightarrow \delta L \Rightarrow \delta \tilde{b}. \tag{1}$$

Due to different strengths of interaction at the different values of the impact parameter, the processes of transition to the spin up and down states will have different probabilities which leads eventually to polarization of $\Lambda$.

In a particular case of $\Lambda$–polarization the relevant transitions of constituent quark $U$ (cf. Fig. 1) will be correlated with the shifts $\delta \tilde{b}$ in impact parameter $\tilde{b}$ of the final $\Lambda$-hyperon, i.e.: $U_\uparrow \rightarrow K^+ + S_\downarrow \Rightarrow -\delta \tilde{b}$ $U_\downarrow \rightarrow K^+ + S_\uparrow \Rightarrow +\delta \tilde{b}. \tag{2}$

Eqs. (2) clarify mechanism of the filtering of spin states: when shift in impact parameter is $-\delta \tilde{b}$ the interaction is stronger compared to the case when shift is $+\delta \tilde{b}$, and the final $S$-quark (and $\Lambda$-hyperon) is polarized negatively.
Figure 1: Transition of the spin-up constituent quark $U$ to the spin-down strange quark.

It is important to note here that the shift of $\tilde{b}$ (the impact parameter of final hyperon) is translated to the shift of the impact parameter of the initial particles according to the relation between impact parameters in the multiparticle production\cite{26}:

$$b = \sum_i x_i \tilde{b}_i. \quad (3)$$

The variable $\tilde{b}$ is conjugated to the transverse momentum of $\Lambda$, but relations between functions depending on the impact parameters $\tilde{b}_i$, which will be used further for the calculation of polarization, are nonlinear and therefore we will use semiclassical correspondence between small and large values of transverse momentum and impact parameter:

small $\tilde{b} \Leftrightarrow$ large $p_\perp$, \quad (4)

large $\tilde{b} \Leftrightarrow$ small $p_\perp$. \quad (5)

We consider production of $\Lambda$ in the fragmentation region, i.e. at large $x_F$ and therefore use approximate relation

$$b \simeq x_F \tilde{b}, \quad (6)$$

which results from Eq. (6)\textsuperscript{3}.

The mechanism of the polarization generation is quite natural and has an optical analogy with the passing of the unpolarized light through the glass of polaroid. The particular mechanism of filtering of spin states is related to the emission of Goldstone bosons by constituent quarks. This picture is more physically apparent and justified than the polarization generation due to multiple scattering of constituent quarks in the effective field developed earlier by the authors in \cite{27}.

We will now obtain an expression for the polarization which takes into account unitarity in the direct channel of reaction and apply chiral quark filtering to conclude on polarization dependence on the kinematical variables.

\textsuperscript{3}We make here an additional assumption on the small values of Feynman $x$ for other particles
2 Polarization dependence on kinematical variables

We use the explicit formulas for inclusive cross–sections of the process

\[ h_1 + h_2 \rightarrow h_3^{\uparrow} + X, \]

where hadron \( h_3 \) is a hyperon whose transverse polarization is measured, obtained in [22]. The main feature of this formalism is an account of unitarity in the direct channel of reaction. The corresponding formulas have the form

\[ d\sigma^{\uparrow,\downarrow}/d\xi = 8\pi \int_0^\infty db I^{\uparrow,\downarrow}(s, b, \xi)/|1 - iU(s, b)|^2, \tag{7} \]

where \( b \) is the impact parameter of the initial particles. Here the function \( U(s, b) \) is the generalized reaction matrix (for unpolarized scattering) which is determined by the basic dynamics of elastic scattering. The elastic scattering amplitude in the impact parameter representation \( F(s, b) \) is related [23, 24, 25] then to the function \( U(s, b) \) by the relation:

\[ F(s, b) = U(s, b)/[1 - iU(s, b)]. \tag{8} \]

This relation allows one to obey unitarity provided inequality \( \text{Im} U(s, b) \geq 0 \) is fulfilled.

The functions \( I^{\uparrow,\downarrow} \) in Eq. (7) are related to the functions \(|U_n|^2\), where \( U_n \) are the multiparticle analogs of the \( U \) [22]. The kinematical variables \( \xi (x_F \text{ and } p_{\perp}) \) describe the state of the produced particle \( h_3 \). Arrows \( \uparrow \) and \( \downarrow \) denote transverse spin directions of the final hyperon \( h_3 \).

Polarization

\[ P = \left\{ \frac{d\sigma^\uparrow}{d\xi} - \frac{d\sigma^\downarrow}{d\xi}\right\}/\left\{ \frac{d\sigma^\uparrow}{d\xi} + \frac{d\sigma^\downarrow}{d\xi}\right\} \]

can be expressed in terms of the functions \( I_-, I_0 \) and \( U \):

\[ P(s, \xi) = \frac{\int_0^\infty db I_-(s, b, \xi)}{2\int_0^\infty db I_0(s, b, \xi)}|1 - iU(s, b)|^2, \tag{9} \]

where \( I_0 = 1/2(I^\uparrow + I^\downarrow) \) and \( I_- = (I^\uparrow - I^\downarrow) \).

Now we turn to the functions \( I^\uparrow \) and \( I^\downarrow \) and will apply chiral quark model to get an information on the hyperon polarization and its dependence on energy, \( x_F \) and \( p_{\perp} \). On the basis of the described chiral quark filtering mechanism we can assume that the functions \( I^\uparrow(s, b, \xi) \) and \( I^\downarrow(s, b, \xi) \) are related to the functions \( I_0(s, b, \xi)|_{\delta b} \) and \( I_0(s, b, \xi)|_{\delta b} \), respectively, i.e.

\[ I_-(s, b, \xi) = I_0(s, b, \xi)|_{\delta b} - I_0(s, b, \xi)|_{\delta b} = 2\delta I_0(s, b, \xi)_{\delta b}. \tag{10} \]
We can connect \( \delta \tilde{b} \) with the radius of quark interaction \( r_{U \rightarrow S}^{\text{flip}} \) responsible for the transition \( U \uparrow \rightarrow S \downarrow \) changing quark spin and flavor:

\[
\delta \tilde{b} \simeq r_{U \rightarrow S}^{\text{flip}}.
\]

Using the above formulas and, in particular, relation (6), we can write the following expression for polarization \( P_\Lambda(s, \xi) \)

\[
P_\Lambda(s, \xi) \simeq x_F r_{U \rightarrow S}^{\text{flip}} \frac{\int_0^\infty b d b I_0'(s, b, \xi) \frac{1}{|1 - i U(s, b)|^2}}{\int_0^\infty b d b I_0(s, b, \xi) / |1 - i U(s, b)|^2},
\]

where \( I_0'(s, b, \xi) = d I_0(s, b, \xi) / d b \). We have made replacement in (11) according to relation (6):

\[
\delta I_0(s, b, \xi) / \delta \tilde{b} \rightarrow d I_0(s, b, \xi) / d b.
\]

It is clear that polarization of \( \Lambda \)-hyperon \( (11) \) should be negative because \( I_0'(s, b, \xi) < 0 \).

The function \( U(s, b) \) is chosen as a product of the averaged quark amplitudes in accordance with the quasi-independence of valence constituent quark scattering in the mean field \[20\].

The generalized reaction matrix \( U(s, b) \) (in a pure imaginary case, which we consider here for simplicity) is the following 4

\[
U(s, b) = i \tilde{U}(s, b) = i g(s) \exp(-M b / \zeta),
\]

where

\[
g(s) \equiv g_0 g_Q^N(s) \equiv g_0 \left[ 1 + \alpha \sqrt{s} / m_Q \right]^N,
\]

\( M \) is the total mass of \( N \) constituent quarks with mass \( m_Q \) in the initial hadrons; \( \alpha \) and \( g_0 \) are the parameters of model. Parameter \( \zeta \) is the one which determines a universal scale for the quark interaction radius, i.e. \( r_Q = \zeta / m_Q \).

Performing integration by parts we can rewrite the expression for polarization \( P_\Lambda(s, \xi) \) in the form:

\[
P_\Lambda(s, \xi) \simeq -x_F r_{U \rightarrow S}^{\text{flip}} \frac{M \int_0^\infty b d b I_0(s, b, \xi) \tilde{U}(s, b) / [1 + \tilde{U}(s, b)]^3}{\zeta \int_0^\infty b d b I_0(s, b, \xi) / [1 + \tilde{U}(s, b)]^2}.
\]

To evaluate polarization dependence on \( x_F \) and \( p_\perp \) we use semiclassical correspondence between transverse momentum and impact parameter values, i.e. (4) and (5).

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4This form leads, in particular, to asymptotic \( \ln^2 s \) dependence for total cross–sections and \( (1/s)^{N+3} f(\theta) \) dependence of differential cross–sections at large angles \[20\].
Choosing the region of small $p_\perp$ we select the large values of impact parameter and therefore we have

$$P_\Lambda(s, \xi) \propto -x_F r^{\text{flip}} u_{S \rightarrow S}^\Lambda \int_{b > R(s)} b db I_0(s, b, \xi) \tilde{U}(s, b),$$  \hspace{1cm} (14)$$

where $R(s) \propto \ln s$ is the hadron interaction radius, which serve as a scale of large and small impact parameter values. At large values of impact parameter $b$: $\tilde{U}(s, b) \ll 1$ for $b \gg R(s)$ and therefore we will have small polarization $P_\Lambda \approx 0$ in the region of small and moderate $p_\perp \leq 1$ GeV/c.

But at small values of $b$ (and large $p_\perp$): $\tilde{U}(s, b) \gg 1$ and the following approximate relations are valid

$$\int_{b < R(s)} b db I_0(s, b, \xi) \tilde{U}(s, b) \left[1 + \tilde{U}(s, b)\right]^3 \simeq \int_{b < R(s)} b db I_0(s, b, \xi) \tilde{U}(s, b)^{-2},$$  \hspace{1cm} (15)$$

since we can neglect unity in the denominators of the integrands. Thus, in (13)

Figure 2: $x_F$ (left panel) and $p_T$ (right panel) dependencies of the $\Lambda$-hyperon polarization

the ratio of two integrals is of order of unity and therefore the energy and $p_\perp$-independent behavior of polarization $P_\Lambda$ takes place at large values of $p_\perp$:

$$P_\Lambda(s, \xi) \propto -x_F r^{\text{flip}} u_{S \rightarrow S}^\Lambda M/\zeta.$$  \hspace{1cm} (16)$$

This flat transverse momentum dependence results from the similar rescattering effects for the different spin states, i.e. spin–flip and non spin-flip interactions undergo similar absorption at short distances and the relative magnitude of this absorption does not depend on energy. It is one of the manifestations of unitarity. The numeric value of polarization $P_\Lambda$ can be large: there are no small factors in (16). In (16) $M$ is proportional to two nucleon masses, the value of parameter

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\( \zeta \simeq 2 \). We expect that \( r_{U \rightarrow S}^{flip} \simeq 0.1 - 0.2 \) fm on the basis of the model \[20\] \[22\], however, this is a crude estimate. The above qualitative features of polarization dependence on \( x_F, p_\perp \) and energy are in a good agreement with the experimentally observed trends \[11\]. For example, Fig. 2 demonstrates that the linear \( x_F \) dependence is in a good agreement with the experimental data in the fragmentation region \( (x_F \geq 0.4) \) where the model should work. Of course, the conclusion on the \( p_\perp \)–independence of polarization is a rather approximate one and deviation from such behavior cannot be excluded.

## 3 Inclusive cross-sections of the unpolarized \( \Lambda \) production

To demonstrate the model self-consistency we consider in this section the unpolarized cross-section of \( \Lambda \)-production:

\[
\frac{d\sigma}{d\xi} = 8\pi \int_0^\infty b db \frac{I_0(s, b, \xi)}{|1 - iU(s, b)|^2}.
\]

(17)

In the fragmentation region we can simplify the problem and consider the process of \( \Lambda \)-production as a quasi two-particle reaction, where the second final particle has a mass \( M^2 \simeq (1 - x_F)s \). The amplitude of this quasi two-particle reaction in the pure imaginary case (which we consider for simplicity) can be written in the form

\[
F(s, p_\perp, x_F) = \frac{is}{x_F\pi^2} \int_0^\infty b db J_0(bp_\perp/x_F) \frac{I_0^{1/2}(s, b, x_F)}{1 + U(s, b)}. \]

(18)

To obtain Eq. 18 we used relations \( b \simeq x_F \tilde{b} \) and due to the fact that the functions \( I_0 \) is quadratic on the the multiparticle analog of the generalized reaction matrix \( U \), we use the relation

\[
I_0^{1/2}(s, b, p_\perp, x_F) = \frac{s}{\pi^2} \int_0^\infty I_0^{1/2}(s, b, \tilde{b}, x_F) J_0(bp_\perp) \tilde{b} db. \]

(19)

Since in the model constituent quarks are considered to form a \( SU(6) \) wave function, \( I_0 = I_{U \rightarrow S}^{1/2} \). The function \( I_{U \rightarrow S}^{1/2}(s, b, x_F) \) according to quasi-independent nature of constituent quark-scattering can be represented then as a product

\[
I_{U \rightarrow S}^{1/2}(s, b, x_F) = \prod_{Q=1}^{N-1} \langle f_Q(s, b) \rangle \langle f_{U \rightarrow S}(s, b, x_F) \rangle,
\]

(20)

where \( N \) is the total number of quarks in the colliding hadrons.
In the model the $b$--dependencies of the amplitudes $\langle f_Q \rangle$ and $\langle f_{U \rightarrow S} \rangle$ are related to the strong form factor of the constituent quark and transitional spin-flip form factor respectively. The strong interaction radius of constituent quark is determined by its mass. We suppose that the corresponding radius of transitional form factor is determined by the average mass $\tilde{m}_Q = (m_U + m_S)/2$ and factor $\kappa < 1$ (which takes into account reduction of the radius due to spin flip) $r_{U \rightarrow S}^{flip} = \kappa \zeta / \tilde{m}_Q$ and the corresponding function $f_{U \rightarrow S}(s, b, x_F)$ has the form

$$f_{U \rightarrow S}(s, b, x_F) = g_{flip}(x_F) \exp \left( -\frac{\tilde{m}_Q}{\kappa \zeta} b \right)$$  \hspace{1cm} (21)$$

The expression for $I_0(s, b, x_F)$ can be rewritten then in the following form:

$$I_0(s, b, x_F) = \bar{g}(x_F) U(s, b) \exp[-\Delta m_Q b / \zeta],$$  \hspace{1cm} (22)$$

where the mass difference $\Delta m_Q \equiv \tilde{m}_Q / \kappa - m_Q$ and $\bar{g}(x_F)$ is the function whose dependence on Feynman $x_F$ in the model is not fixed.

Now we can consider $p_\perp$- and $x_F$-dependencies of the $\Lambda$-hyperon production cross-section and we start with angular distribution\(^5\). The corresponding amplitude $F(s, p_\perp, x_F)$ can be calculated analytically. To do so we continue the amplitudes $F(s, \beta, x_F)$, $\beta = b^2$, where

$$F(s, \beta, x_F) = \frac{1}{x_F^2} \frac{I_0^{1/2}(s, \beta, x_F)}{1 + U(s, \beta)}$$

to the complex $\beta$–plane and transform the Fourier–Bessel integral over impact parameter into the integral in the complex $\beta$–plane over the contour $C$ which goes around the positive semiaxis. Then for the amplitude $F(s, p_\perp, x_F)$ the following representation takes place:

$$F(s, p_\perp, x_F) = -\frac{is}{2\pi^2} \int_C d\beta F(s, \beta, x_F) K_0(\sqrt{-p_\perp^2 \beta / x_F})$$  \hspace{1cm} (23)$$

where $K_0(x)$ is the modified Bessel function. The amplitude $F(s, \beta, x_F)$ has the poles in the $\beta$–plane determined by equation

$$1 + U(s, \beta) = 0.$$  \hspace{1cm} (24)$$

The solutions of this equation can be written as

$$\beta_n(s) = \frac{\zeta^2}{M^2} \left\{ \ln g(s) + i\pi n \right\}, \quad n = \pm 1, \pm 3, \ldots \hspace{1cm} (25)$$

One should remember that all formulas and figures below are valid for the fragmentation region only, i.e. for $x_F > 0.4$.  

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The amplitude \( F(s, \beta, x_F) \) besides the poles has a branching point at \( \beta = 0 \). Therefore the amplitude \( F(s, p_{\perp}, x_F) \) can be represented as a sum of the poles contribution and the contribution of the cut:

\[
F(s, p_{\perp}, x_F) = F_p(s, p_{\perp}, x_F) + F_c(s, p_{\perp}, x_F)  \tag{26}
\]

The poles contribution has an exponential dependence on \( p_{\perp} \)

\[
F_p(s, p_{\perp}, x_F) \simeq isG_p(s, x_F) \sum_{k=1}^{\infty} \tau^k(-p_{\perp}/x_F)\varphi_k[R(s), p_{\perp}],  \tag{27}
\]

where the parameter \( \tau(-p_{\perp}/x_F) \)

\[
\tau(-p_{\perp}/x_F) = \exp(-\frac{\pi \zeta p_{\perp}}{x_F})
\]

the function \( G_p(s, x_F) \) is

\[
G_p(s, x_F) = \left[ \frac{\bar{g}(x_F)}{x_F^4g(s)} \right]^{1/2} [g(s)]^{-\frac{\Delta m_Q}{2\bar{M}}} 
\]

and \( \varphi_k[R(s), p_{\perp}] \) is the oscillating functions of \( p_{\perp} \), the hadron interaction radius \( R(s) \) determines the period of these oscillations and it has slow energy dependence like \( \ln^{1/2} s \).

The cut contributions has power-like dependence on \( p_{\perp} \)

\[
F_c(s, p_{\perp}, x_F) \simeq isG_c(s, x_F)(1 + \frac{p_{\perp}^2}{x_F^2\bar{M}^2})^{-3/2},  \tag{28}
\]

where \( \bar{M} = (M - \Delta m_Q)/2\zeta \) and the function \( G_c(s, x_F) \) has the form

\[
G_p(s, x_F) = \left[ \frac{\bar{g}(x_F)}{x_F^4g(s)} \right]^{1/2} [g(s)]^{-\frac{1}{2}} 
\]

Calculation of poles and cut contributions are similar to the case of elastic scattering \([28]\).

The poles and cut contributions determine the behaviour of the inclusive cross-section of \( \Lambda \) production at moderate and large values of \( p_{\perp} \) correspondingly, i.e. it will have in the region of large \( p_{\perp} \) power-like dependence on \( p_{\perp} \):

\[
\frac{d\sigma}{d\xi} \propto G^2_c(s, x_F)(1 + \frac{p_{\perp}^2}{x_F^2\bar{M}^2})^{-3},  \tag{29}
\]
while at smaller values of $p_{\perp}$ inclusive cross-section would have the exponential $p_{\perp}$-dependence:

$$\frac{d\sigma}{d\xi} \propto G_p^2(s, x_F) \exp\left(-\frac{2\pi \zeta p_{\perp}}{M x_F}\right). \tag{30}$$

The data for the $\Lambda$-hyperon production are available at the moderate values of $p_{\perp}$ and the experimental fits to the data \cite{29} of the form

$$A(1 - x_F)^n e^{-B(x_F)p_{\perp}}$$

just follow to Eq. (30) when relevant parameterization for the function $\bar{g}(x_F)$ is chosen. At high values of $p_{\perp}$ power-like dependence should take place according to Eq. \cite{29}. In the energy region of $\sqrt{s} \lesssim 2$ TeV the functions $G_p$ and $G_c$ have very slow variation with energy due to the numerical values of parameters \cite{30}.

### 4 Comments on SSA in pion production processes

SSA is an interesting topic not only in the field of hyperon polarization. The new experimental expectations are related to the experiments at RHIC with polarized proton beams and new experimental data obtained by the STAR collaboration have already demonstrated significant spin asymmetry in the $\pi^0$–production similar to the one observed in the fragmentation region at FNAL \cite{31, 32}.

We would like to make a brief comment on this subject and to note that the reverse to the filtering mechanism can be used for the explanation of the SSA in pion production observed at FNAL and recently at RHIC in the fragmentation region. In the initial state of these reaction the proton is polarized and can be represented in the simple SU(6) model as following:

$$p_\uparrow = \frac{5}{3} U_\uparrow + \frac{1}{3} U_\downarrow + \frac{1}{3} D_\uparrow + \frac{2}{3} D_\downarrow. \tag{31}$$

The relevant process for $\pi^+$–production is

$$U_\uparrow \rightarrow \pi^+ + D_\downarrow,$$

which leads to a negative shift in the impact parameter and consequently to the positive asymmetry $A_N$, while the corresponding process for the $\pi^-$–production

$$D_\downarrow \rightarrow \pi^- + U_\uparrow.$$

It leads to the positive shift in impact parameter and respectively to the negative asymmetry $A_N$. Asymmetry $A_N$ in the fragmentation region should have similar to polarization linear $x_F$–dependence which is in agreement with the observed
Figure 3: \( x_F \)-dependence of the asymmetry \( A_N \) in \( \pi^0 \)-production (filled circles–E704 data \([31]\), empty circles–STAR data \([32]\)).

FNAL and RHIC experimental data \([31,32]\) at \( x_F > 0.4 \). As for the neutral \( \pi^0 \)-production the combination of \( U \) and \( D \)-quarks with up and down polarization makes contributions to cross-sections and asymmetry. On the basis of the simple SU(6) model we can assume that the \( U \)-quark with up polarization would contribute mainly in the fragmentation region. Then the \( \pi^0 \)-production should have positive asymmetry. The corresponding behavior and experimental data obtained at FNAL and RHIC are depicted on Fig. 3. Linear \( x_F \)-dependence agrees with the experimental data at large \( x_F \) (fragmentation region, \( x_F > 0.4 \)).

The expression for the unpolarized inclusive cross-section of the pion production in the fragmentation region can be obtained in a similar way to the corresponding cross-section of \( \Lambda \) production. It should be noted that Eq. (29) leads to \( p_{\perp}^{-6} \) dependence at large \( p_{\perp} \) and is valid also for the pion production. It is in agreement with \( p_{\perp}^{-N} \) (with the exponent \( N = 6.2 \pm 0.6 \)) dependence of the inclusive cross-section of \( \pi^0 \)-production observed in forward region at large \( p_{\perp} \) at RHIC \([33]\). The exponent \( N \) does not depend on \( x_F \) and choosing relevant function \( \bar{g}(x_F) \) the \( (1 - x_F)^{5.1 \pm 0.6} \) dependence of experimental data can be reproduced.

**Conclusion**

The proposed mechanism deals with effective degrees of freedom and takes into account collective aspects of QCD dynamics. Together with unitarity, which is an essential ingredient of this approach, it allows to obtained results for polarization dependence on kinematical variables in agreement with the experimental behavior of \( \Lambda \)-hyperon polarization, i.e. linear dependence on \( x_F \) and flat dependence on \( p_{\perp} \) at large \( p_{\perp} \) in the fragmentation region are reproduced. Those dependencies together with the energy independent behavior of polarization at large transverse momenta are the straightforward consequences of this model.
We discuss here particle production in the fragmentation region. In the central region where correlations between impact parameter of the initial and impact parameters of the final particles being weakened, the polarization cannot be generated due to chiral quark filtering mechanism. Moreover, it is clear that since antiquarks are produced through spin-zero Goldstone bosons we should expect $P_A \simeq 0$. The chiral quark filtering is also relatively suppressed when compared to direct elastic scattering of quarks in effective field and therefore should not play a role in the reaction $pp \rightarrow pX$ in the fragmentation region, i.e. protons should be produced unpolarized. These features take place in the experimental data set.

The application of this mechanism to description of polarization of other hyperons is more complicated problem, since they could have two or three strange quarks and spins of $U$ and $D$ quarks can also make contributions into their polarizations.

Finally, it was shown that the mechanism reversed to chiral quark filtering can provide description of the SSA in $\pi^0$ production measured at FNAL and recently at RHIC in the fragmentation region and it leads to the energy independence of the asymmetry.

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