ФИЗИКА ЭЛЕМЕНТАРНЫХ ЧАСТИЦ И АТОМНОГО ЯДРА
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We discuss a nonperturbative mechanism for generation of the single-spin asymmetries in hadron interactions. It is based on the chiral quark model combined with unitarity and impact parameter picture and provides explanation for the experimental regularities observed under the measurements of the spin asymmetries.

Обсуждается механизм возникновения односпиновых асимметрий в адронных процессах, основанный на использовании киральной кварковой модели совместно с учетом унитарности и представления прицельного параметра. Механизм позволяет дать объяснение наблюдаемым экспериментально зависимостям в поведении этих асимметрий.

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INTRODUCTION

Studies of the single-spin asymmetries (SSAs) is a sensitive tool to probe QCD at small and large distances. Experimentally, significant SSAs were observed in various processes of elastic scattering, inclusive and exclusive hadron production.

The processes of hadron interactions are complicated, there is no proof of factorization theorem for these processes and it could result from the real absence of hard and soft parts of interaction factorization in hadronic reactions. The origin of SSA in these reactions is not clear. Despite significant efforts in theoretical studies devoted to this problem, the phenomenological success is rather limited; at the moment there is no comprehensive approach able to describe the existing set of experimental data on polarization, asymmetries, spin correlations and the unpolarized cross sections. Theoretically, there are various approaches to generation of the nonzero SSA but prevailing number of studies of the SSAs in the field are based on *assumed* extended factorization in perturbative QCD with considerations of the Sivers (structure functions) and/or Collins (fragmentation functions) mechanisms [1–5] combined sometime with inclusion of the higher twists contributions to the scattering amplitude of the seemingly to be hard parton subprocess [6–8].

The decreasing dependence of SSA with p_T — common feature for the listed above approaches — has not been observed experimentally. The experimental data including the most recent data obtained at RHIC [9], are consistent with a flat transverse momentum dependence at $p_T \geqslant 1$ GeV/c. Another important point is related to the unpolarized inclusive cross section. For example, it has also been demonstrated [10] that the description of the inclusive cross section for π^0 production, at the energies lower than the RHIC energies, also meets difficulties in the framework of the perturbative QCD. Deviation from the pQCD scaling is mostly noticeable in the forward region where the most significant asymmetry in the π^0 production in $pp_\uparrow \to \pi^0 X$ has also been observed by STAR Collaboration at RHIC [9] at $\sqrt{s} = 200$ GeV (in the fragmentation region of the polarized proton).

Of course, more experimental data are needed to perform a conclusive test of various theoretical approaches, and their predictions should be more specified and elaborated for the observables at the hadronic level. In this connection it should be noted that one of the most interesting and persistent spin phenomena is a very significant polarization of Λ hyperons which has been discovered almost three decades ago in collisions of unpolarized hadron beams [12]. It should be also noted that the asymmetry $A_N = 0$ in the neutral pion production in the backward and midrapidity regions [13,14]. SSA has also zero value in the $pp_\uparrow \to pX$, while $A_N \neq 0$ in the $pp_\uparrow \to nX$ [15] in the polarized proton fragmentation region. The approaches based on the assumed pQCD factorization meet in these processes the problems mentioned above.

Thus, it is (depending more or less on the particular personal taste) evident that the problems mentioned above can be related to the illegitimate use of the methods based on perturbative expansion, factorization and accounts for higher twists in the region and in the processes where they actually cannot be valid, and it is the kinematical region of the modern experiments dealing with rather modest transverse momenta and energies. In contrast, it might happen that SSAs originate from the genuine nonperturbative sector of QCD (cf., e.g., [16]). Such point of view, i.e., that the polarization has its roots in the nonperturbative sector of QCD, was widely used in the earlier models and becomes less isolated one nowadays. *In the nonperturbative sector of QCD*, the two phenomena, confinement and chiral symmetry spontaneous breaking (χSB) [17], should be reproduced. The relevant scales are characterized by the parameters $\Lambda_{\rm QCD}$ and Λ_χ, respectively. Chiral $SU(3)_L \times SU(3)_R$ symmetry is spontaneously broken at the distances between these two scales. The χSB leads to generation of quark masses and appearance of quark condensates. It describes transition of current into constituent quarks. Constituent quarks in its turn are quasiparticles, i.e., they are a coherent superposition of bare quarks and their masses are 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 a picture provides reasonable values for the static characteristics of hadrons, for instance, their magnetic moments. The other well-known result is the appearance of the Goldstone bosons. It has been successfully applied for the explanation of the nucleon spin structure [19] including the most recent results obtained at JLab [20].

It is necessary to note that the *structure functions* are represented by the distorted parton distributions in the impact parameter plane in the polarized case [3]. In this work the approach based on nonperturbative QCD has been used to relate Λ polarization with large magnitude of the transverse flavor dipole moment of the transversely polarized baryons.

The instanton-induced SSA generation [21, 22] relates those asymmetries to a genuine nonperturbative QCD interaction. It should be noted that the physics of instantons (cf., e.g., [23]) can provide microscopic explanation for the χSB*.

We discuss here the SSA generation based on chiral quark model ideas (cf., e.g., [17]) and the filtering spin states related to the account of unitarity in the s channel. It connects polarization with asymmetry in the position (impact parameter) space. We show that the common features of SSA measurements at RHIC and Tevatron (linear increase of asymmetry with x_F and flat transverse momentum dependence at $p_T > 1$ GeV/c) can be reproduced and described in the framework of the semiclassical picture based on the further development of the chiral quark model suggested in [24] and results of its adaptation for the treatment of the polarized and unpolarized inclusive cross sections. The above-mentioned data obtained at RHIC [11] for the unpolarized inclusive cross section can be simultaneously described. Consistency with other new experimental regularities found at RHIC are discussed as well. Preliminary version of this paper can be found in [18].

1. SEMICLASSICAL MECHANISM OF SSA GENERATION

As was argued, the SSA could originate from the nonperturbative QCD and is related to the mechanism of spontaneous chiral symmetry breaking (χSB) in QCD [25], leading to generation of quark masses and appearance of quark condensates.

Thus we consider a hadron as an extended object consisting of the valence constituent quarks located in the central core which is embedded into a quark condensate. Collective excitations of the condensate are the Goldstone bosons,

*We are grateful to Dmitri Diakonov for the interesting communication on this matter regarding the polarization phenomena.

and the constituent quarks interact via exchange of Goldstone bosons [26]. This interaction is mainly due to a pion field and has therefore a spin-flip nature.

At the first stage of hadron interaction common effective self-consistent field appears. This field is generated by $\bar{Q}Q$ pairs and pions interacting with quarks. The time of generation of the effective field $t_{\rm eff}$:

$$t_{\rm eff} \ll t_{\rm int},$$

where $t_{\rm int}$ is the total interaction time. This assumption on the almost instantaneous generation of the effective field obtained some support in the very short thermalization time revealed in heavy-ion collisions at RHIC [27].

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 used in the model [24] which has been applied to description of elastic scattering and hadron production [28].

In the initial state of the reaction $pp_\uparrow \to \pi^0 X$ the proton is polarized and can be represented in the simple $SU(6)$ model as following:

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

We exploit the common feature of chiral quark models; namely, the constituent quark $Q_\uparrow$ with transverse spin in up-direction can fluctuate into Goldstone boson and into another constituent quark $Q'_\downarrow$ with opposite spin direction, i.e., perform a spin-flip transition [19]:

$$Q_\uparrow \to {\rm GB} + Q'_\downarrow. \tag{2}$$

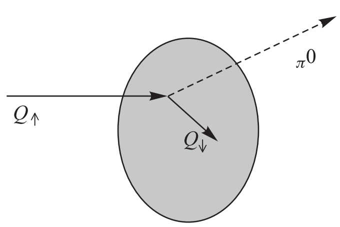

Fig. 1. Schematic view of π^0 production in polarized proton–proton interaction

The π^0 fluctuations of quarks do not change the quark flavor, and assuming they have equal probabilities in the processes:

$$U_{\uparrow,\downarrow} \to \pi^0 + U_{\downarrow,\uparrow} \quad \text{and} \quad D_{\uparrow,\downarrow} \to \pi^0 + D_{\downarrow,\uparrow}, \tag{3}$$

the production of π^0 by the polarized proton $p_\uparrow$ in this simple $SU(6)$ picture can be regarded as a result of the fluctuation of the constituent quark $Q_\uparrow$ ($Q = U$ or D) in the effective field into the system $\pi^0 + Q_\downarrow$ (Fig. 1).

The contributions to the cross-sections difference of the quarks polarized in opposite directions compensate each other (as it will be clear in what follows), and it is not the case for the π^0 production in the unpolarized case. Therefore, the asymmetry A_N should obey the inequality $|A_N(\pi^0)| \leqslant 1/3$.

To compensate quark spin flip $\delta\mathbf{S}$, an orbital angular momentum $\delta\mathbf{L} = -\delta\mathbf{S}$ should be attributed to the final state of reaction (2). The presence of $\delta\mathbf{L}$ in its turn means a shift in the impact parameter value of the Goldstone boson π^0:

$$\delta\mathbf{S} \Rightarrow \delta\mathbf{L} \Rightarrow \delta\tilde{\mathbf{b}}.$$

Due to different strengths of interaction at the different impact distances, i.e.,

$$\begin{aligned} p_\uparrow &\Rightarrow Q_\uparrow \to \pi^0 + Q_\downarrow \Rightarrow -\delta\tilde{\mathbf{b}}, \\ p_\downarrow &\Rightarrow Q_\downarrow \to \pi^0 + Q_\uparrow \Rightarrow +\delta\tilde{\mathbf{b}}, \end{aligned} \tag{4}$$

the processes of transition $Q_\uparrow$ and $Q_\downarrow$ to π^0 will have different probabilities. It eventually leads to nonzero asymmetry $A_N(\pi^0)$. Equations (4) clarify mechanism of the SSA generation: when shift in impact parameter is $-\delta\tilde{\mathbf{b}}$, the interaction is stronger than when the shift is $+\delta\tilde{\mathbf{b}}$, and the asymmetry $A_N(\pi^0)$ is positive. It is important to note here that the shift in $\tilde{\mathbf{b}}$ (the impact parameter of final pion) is equivalent to the shift of the impact parameter of the initial proton according to the relation between impact parameters in the multiparticle production [29]:

$$\mathbf{b} = \sum_i x_i \tilde{\mathbf{b}}_i. \tag{5}$$

The variable $\tilde{b}$ is conjugated to the transverse momentum of π^0, but relations between functions depending on the impact parameters $\tilde{b}_i$, which will be used further for the calculation of asymmetry, are nonlinear and therefore we are using the semiclassical correspondence between small and large values of transverse momentum and impact parameter:

$$\text{small } \tilde{b} \Leftrightarrow \text{large } p_T \quad \text{and} \quad \text{large } \tilde{b} \Leftrightarrow \text{small } p_T. \tag{6}$$

We consider production of π^0 in the fragmentation region, i.e., at large x_F and therefore use the approximate relation

$$b \simeq x_F \tilde{b}, \tag{7}$$

which results from Eq. (5) with additional assumption on the small values of Feynman x_F for the other particles. In the symmetrical case of pp interactions the model assumes equal average multiplicities in the forward and backward hemispheres. It supposes also small momentum transfer between the two sides. This is based on the arguments by Chou and Yang [30].

2. ASYMMETRY AND INCLUSIVE CROSS SECTION

We apply chiral quark semiclassical mechanism which takes into account unitarity in the direct channel to obtain qualitative conclusions on asymmetry dependence on the kinematical variables.

The main feature of the mechanism is an account of unitarity in the direct channel of reaction. The corresponding formulas for inclusive cross sections of the process

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

where hadron h_3 in this particular case is π^0 meson and h_1, h_2 are protons, were obtained in [31] and have the following form:

$$\frac{d\sigma^{\uparrow,\downarrow}}{d\xi} = 8\pi \int\limits_0^\infty b\, db\, I^{\uparrow,\downarrow}(s,b,\xi)/|1 - iU(s,b)|^2, \tag{8}$$

where b is the impact parameter of the initial protons. Here the function $U(s,b)$ is the generalized reaction matrix (averaged over initial spin states) which is determined by the basic dynamics of the elastic scattering. The elastic scattering amplitude in the impact parameter representation $F(s,b)$ is then given [32] by the relation:

$$F(s,b) = \frac{U(s,b)}{1 - iU(s,b)}. \tag{9}$$

This equation allows one to obey unitarity provided inequality $\operatorname{Im} U(s,\, b) \geqslant 0$ is fulfilled. The functions $I^{\uparrow,\downarrow}$ in Eq. (8) are related to the functions $U_n^{\uparrow,\downarrow}$ — the multiparticle analogs of the function U [31] in the polarized case. The kinematical variables ξ (x_F and p_T, for example) describe the state of the produced particle h_3. Arrows $\uparrow$ and $\downarrow$ denote transverse spin directions of the polarized proton h_2.

Asymmetry A_N can be expressed in terms of the functions I_-, I_0, and U:

$$A_N(s,\, \xi) = \frac{\int\limits_0^\infty b\, db\, I_-(s,b,\xi)/|1 - iU(s,b)|^2}{2\int\limits_0^\infty b\, db\, I_0(s,b,\xi)/|1 - iU(s,b)|^2}, \tag{10}$$

where $I_0 = 1/2(I^{\uparrow} + I^{\downarrow})$, $I_- = (I^{\uparrow} - I^{\downarrow})$ and I_0 obey the sum rule

$$\int I_0(s,b,\xi) d\xi = \bar{n}(s,b) \operatorname{Im} U(s,b),$$

here $\bar{n}(s,b)$ stands for the mean multiplicity in the impact parameter representation.

On the basis of the described mechanism we can assume that the functions $I^{\uparrow}(s,b,\xi)$ and $I^{\downarrow}(s,b,\xi)$ are related to the functions $\frac{1}{3}I_0(s,b,\xi)|_{\tilde{b}-\delta\tilde{b}}$ and $\frac{1}{3}I_0(s,b,\xi)|_{\tilde{b}+\delta\tilde{b}}$, respectively, i.e.,

$$I_-(s,b,\xi) = \frac{1}{3}[I_0(s,b,\xi)|_{\tilde{b}-\delta\tilde{b}} - I_0(s,b,\xi)|_{\tilde{b}+\delta\tilde{b}}] = -\frac{2}{3}\frac{\delta I_0(s,b,\xi)}{\delta\tilde{b}}\delta\tilde{b}. \tag{11}$$

We can connect $\delta\tilde{b}$ with the radius of quark interaction $r_Q^{\rm flip}$ responsible for the transition changing quark spin:

$$\delta\tilde{b} \simeq r_Q^{\rm flip}.$$

Using the above relations and, in particular, (7), we can write the following expression for asymmetry $A_N^{\pi^0}$:

$$A_N^{\pi^0}(s,\xi) \simeq -x_F r_Q^{\rm flip} \frac{1}{3} \frac{\int\limits_0^\infty b\,db\, I_0'(s,b,\xi)db/|1-iU(s,b)|^2}{\int\limits_0^\infty b\,db\, I_0(s,b,\xi)/|1-iU(s,b)|^2}, \tag{12}$$

where $I_0'(s,b,\xi) = dI_0(s,b,\xi)/db$. In Eq. (12) we made replacement according to relation (7):

$$\frac{\delta I_0(s,b,\xi)}{\delta\tilde{b}} \Rightarrow \frac{x_F dI_0(s,b,\xi)}{db}.$$

It is clear that $A_N^{\pi^0}(s,\,\xi)$ (12) should be positive 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 self-consistent mean field [24]. The generalized reaction matrix $U(s,\,b)$ (in a pure imaginary case, which we consider here for simplicity) has the following form:

$$U(s,b) = i\tilde{U}(s,b) = ig(s)\exp\left(-\frac{Mb}{\zeta}\right), \tag{13}$$

where the function $g(s)$ increases at large values of s like a power of energy:

$$g(s) = \left[1 + \alpha\frac{\sqrt{s}}{m_Q}\right]^N,$$

M is the total mass of N constituent quarks with mass m_Q in the initial hadrons, and parameter ζ determines a universal scale for the quark interaction radius in the model, i.e., $r_Q = \zeta/m_Q$.

To evaluate asymmetry dependence on x_F and p_T we use semiclassical correspondence between transverse momentum and impact parameter values, Eq. (6). Performing integration by parts we can rewrite the expression for the asymmetry in the form:

$$A_N^{\pi^0}(s,\xi) \simeq x_F r_Q^{\rm flip} \frac{M}{3\zeta} \frac{\int\limits_0^\infty b\,db\, I_0(s,b,\xi)\tilde{U}(s,b)/[1+\tilde{U}(s,b)]^3}{\int\limits_0^\infty b\,db\, I_0(s,b,\xi)/[1+\tilde{U}(s,b)]^2}. \tag{14}$$

At small values of b the values of U matrix are large, and we can neglect unity in the denominators of the integrands (however it is rather rough approximation valid only at high enough energies). Thus the ratio of the two integrals (after integration by parts of nominator in Eq. (14)) is of the order of unity, i.e., the energy and p_T-independent behavior of asymmetry $A_N^{\pi^0}$ takes place at the values of transverse momentum $p_T \gg x_F/R(s)$:

$$A_N^{\pi^0}(s,\,\xi) \sim x_F r_Q^{\rm flip} \frac{M}{3\zeta}. \tag{15}$$

Such a flat transverse momentum dependence of asymmetry results from the similarity of the rescattering effects for the different spin states, i.e., spin-flip and spin-nonflip 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 the unitarity. The numeric value of polarization $A_N^{\pi^0}$ can be significant. Indeed, there is no small factor in (15). In Eq. (15) M is equal to the total mass of the constituent quarks in the colliding nucleons, the value of parameter $\zeta \simeq 2$. We expect that $r_Q^{\rm flip} \sim 0.1$ fm on the basis of the model estimate [24, 31]. The above qualitative features of asymmetry dependence on x_F, p_T, and energy are in agreement with the experimentally observed trends [11]. For example, Fig. 2 demonstrates that the linear x_F and p_T dependences are in agreement with the experimental data of STAR Collaboration at RHIC [11] in the fragmentation region ($x_F \geqslant 0.4$). It is this region where the model should be applicable. Of course, these dependences of polarization are the qualitative ones and deviations cannot be excluded. The same dependences are compared with the FNAL E704 data [33] (Fig. 3). Those dependences are valid in high-energy approximation and therefore have been compared with FNAL and RHIC data only. Nevertheless, they are in a qualitative agreement with the lower energy data also [34]. Comparison with experimental data allows one to estimate the value $r_Q^{\rm flip} \simeq 0.07$ fm.

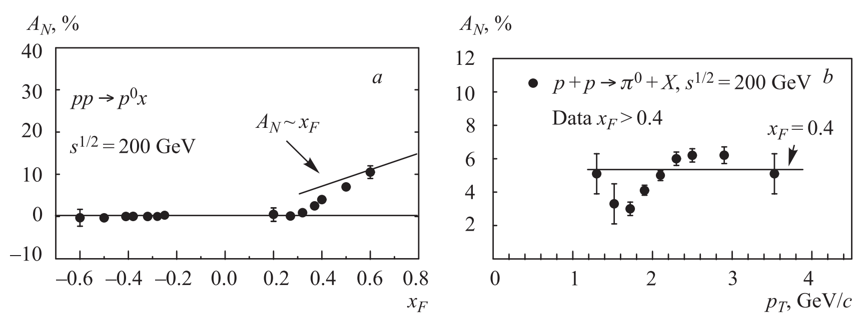

Fig. 2. x_F (*a*) and p_T (*b*) dependences of the asymmetry A_N in the process $p + p_\uparrow \to \pi^0 + X$ at RHIC, experimental data from [9]

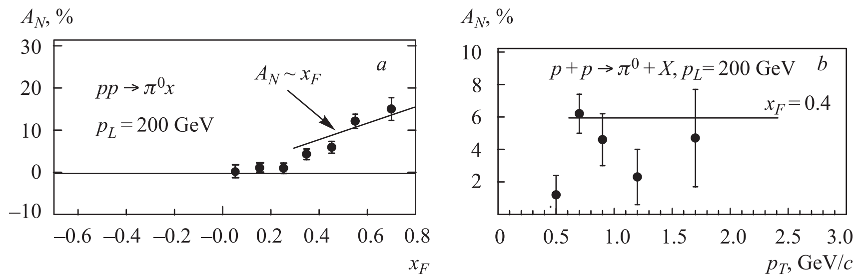

Fig. 3. x_F (*a*) and p_T (*b*) dependences of the asymmetry A_N in the process $p + p_\uparrow \to \pi^0 + X$ at FNAL, experimental data from [33]

Similar mechanism should generate SSA in the production of charged pions. The relevant process for π^+ production in polarized $pp_\uparrow$ interactions

$$U_\uparrow \to \pi^+ + D_\downarrow$$

leads to a negative shift in the impact parameter and consequently to the positive asymmetry A_N, while the corresponding process for the π^- production

$$D_\downarrow \to \pi^- + U_\uparrow$$

leads to the positive shift in impact parameter and, respectively, to the negative asymmetry A_N. Asymmetry A_N in the $\pi^\pm$ production in the fragmentation region of polarized proton should have linear x_F dependence at $x_F > 0.4$ and flat p_T dependence at large p_T. Those dependences are similar to the ones depicted in Fig. 2 for π^0 production. It should be noted that at large transverse momenta, asymmetries are energy-independent at high energies.

Choosing the region of small p_T we select then the large values of impact parameter and obtain

$$A_N^{\pi^0}(s, \xi) \sim x_F r_Q^{\text{flip}} \frac{M}{3\zeta} \frac{\int\limits_{b>R(s)} b\, db\, I_0(s, b, \xi)\tilde{U}(s, b)}{\int\limits_{b>R(s)} b\, db\, I_0(s, b, \xi)}, \tag{16}$$

where $R(s) \sim \ln s$ is the hadron interaction radius, which serves as a scale separating large and small impact parameter regions. Equation (16) does not allow one to draw a definite conclusion on asymmetry behaviour. Its dependence relies on the unknown function $I_0(s, b, \xi)$. Nevertheless, it would be interesting to have at least qualitative estimates of the size of single-spin asymmetries in

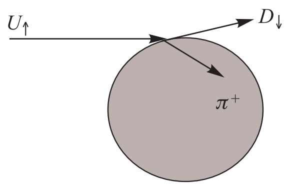

Fig. 4. Schematic view of the constituent quark transition on the boundary of effective field

the small momentum transferred region. It should be noted in this connection, that this region includes the interactions at the boundary of the effective field localization domain (cf. Fig. 4). Therefore, in principle, the asymmetry, which is determined by the variation $\delta I_0(s, b, \xi)/\delta\tilde{b}$, could have significant values due to a large gradient of the interaction intensity in the boundary region. Thus, at the values of transverse momenta corresponding to the constituent quark transition, one can observe very significant asymmetries as it really happens in the forward neutron production at RHIC (cf., e.g., [13]). Unfortunately, we cannot provide the quantitative estimates of the intensity interaction gradient. We can just point out to the possibility, that the similar phenomena should take place in other reactions such as polarization of Λ hyperons, and it would be important therefore to scan experimentally the region of small transverse momenta in the forward production by measurements in narrow bins of transverse momentum. It should be noted in this connection that the chiral quark fluctuation in the effective field with spin flip is relatively suppressed when compared to direct elastic scattering of quarks and therefore does not play a role, e.g., in the reaction $pp_\uparrow \to pX$ in the fragmentation region, but it should not be suppressed in $pp_\uparrow \to nX$. These features really can be observed experimentally: asymmetry A_N is consistent with zero for proton production and significantly deviates from zero for neutron production in the forward region.

To underline the model self-consistency we will demonstrate that it is able to describe the unpolarized cross section of π^0 production also (Sec. 5).

3. SPIN FILTERING AND THE HYPERON POLARIZATION

Chiral quark spin filtering can be used for the explanation of the hyperon polarization [35]. Note that polarization of Λ hyperon has the same generic dependence on x_F and p_T as the asymmetries in the pion production. In this section we consider the origin of the hyperon polarization in the processes where particles in the initial state are unpolarized.

Experimentally, the process of Λ production has been studied more extensively than other hyperon production processes. Observed pattern of the hyperon polarization is well known, being stable for a long time*.

*Polarization of Λ hyperons produced in the unpolarized inclusive pp interactions is negative and energy-independent, it increases linearly with x_F at large transverse momenta ($p_T \geqslant 1$ GeV/c), and for such values of transverse momenta is almost p_T-independent [12].

The main idea is the filtering or discrimination between the two initial spin states of equal probability due to different strength of interactions in the course of scattering in the effective field. The description of spin filtering is considered on the basis of chiral quark model, formulas for inclusive cross section (with the account for the unitarity) [31], and notion on the quasi-independent nature of valence quark scattering in the effective field.

We will use the already discussed 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:

$$Q_\uparrow \to GB + Q'_\downarrow \to Q + \bar{Q}' + Q'_\downarrow. \tag{17}$$

To compensate quark spin flip an orbital angular momentum should be generated in final state of reaction (17). The presence of this orbital momentum $\delta\mathbf{L}$ 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 Λ)

$$\delta\mathbf{S} \Rightarrow \delta\mathbf{L} \Rightarrow \delta\tilde{\mathbf{b}}.$$

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 will lead eventually to polarization of Λ.

In a particular case of Λ polarization, the relevant transitions of constituent quark U (cf. Fig. 5) will be correlated with the shifts $\delta\tilde{b}$ in impact parameter $\tilde{b}$ of the final Λ hyperon, i.e.:

$$\begin{aligned} U_\uparrow &\to K^+ + S_\downarrow \Rightarrow -\delta\tilde{\mathbf{b}}, \\ U_\downarrow &\to K^+ + S_\uparrow \Rightarrow +\delta\tilde{\mathbf{b}}. \end{aligned} \tag{18}$$

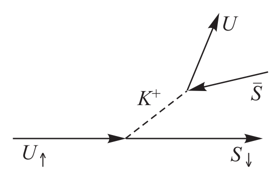

Fig. 5. Transition of the spin-up constituent quark U to the spin-down strange quark

Equations (18) clarify mechanism of the spin-states filtering: when shift in impact parameter is $-\delta\tilde{\mathbf{b}}$, the interaction is stronger compared to the case when shift is $+\delta\tilde{\mathbf{b}}$, and the final S quark (and Λ hyperon) is polarized negatively.

The shift of $\tilde{\mathbf{b}}$ (the impact parameter of final hyperon) is translated then to the shift of the impact parameter of the initial particles.

The mechanism of the polarization generation is quite natural and it has analogy in optics with the passing of the unpolarized light through the glass of polaroid. Spin states filtering is related to emission of Goldstone bosons by constituent quarks.

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

4. Λ-POLARIZATION DEPENDENCE ON KINEMATICAL VARIABLES

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

$$h_1 + h_2 \to h_3^{\uparrow} + X,$$

where hadron h_3 is a hyperon whose transverse polarization is measured, obtained in [31]. Calculation of polarization of Λ proceeds the same steps as those described in Sec. 1, i.e., 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:

$$P(s,\xi) = \frac{\int\limits_0^\infty b\,db\,I_-(s,b,\xi)/|1-iU(s,b)|^2}{2\int\limits_0^\infty b\,db\,I_0(s,b,\xi)/|1-iU(s,b)|^2}, \tag{19}$$

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

We can connect $\delta\tilde{b}$ with the radius of quark interaction $r_{U\to S}^{\rm flip}$ responsible for the transition $U_{\uparrow} \to S_{\downarrow}$ changing quark spin and flavor:

$$\delta\tilde{b} \simeq r_{U\to S}^{\rm flip}.$$

Using the formulas from previous sections, we will arrive to the energy and p_T-independent behavior of polarization P_Λ at small values of b (and large p_T):

$$P_\Lambda(s,\xi) \propto -x_F r_{U\to S}^{\rm flip} \frac{M}{\zeta}. \tag{20}$$

A numeric value of polarization P_Λ can be large: there are again no small factors in (20). The above qualitative features of polarization dependence on x_F, p_T, and energy are in good agreement with the experimentally observed trends [12]. For example, Fig. 6 demonstrates that the linear x_F dependence is in good agreement with the experimental data in the fragmentation region ($x_F \geqslant 0.4$) where the model should work. Of course, the conclusion on the p_T independence of polarization is a rather approximate one and deviation from such behavior cannot be excluded. It should also be noted that the spin filtering mechanism of hyperon polarization does not give contribution to the transverse polarization of Λ in the inclusive reaction of quasireal photoproduction as it is supposed by the experimental data obtained by HERMES Collaboration [36] which reveals positive polarization of Λ and its strikingly different dependence on the light-cone momentum fraction. As was noted in [36], the positive polarization might indicate that the $\gamma \to s\bar{s}$ hadronic component of the photon plays a significant role in the quasireal photoproduction instead.

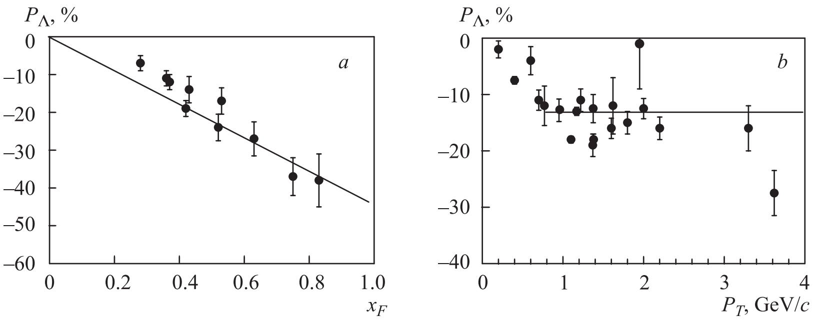

Fig. 6. x_F (*a*) and p_T (*b*) dependences of the Λ-hyperon polarization

At small transverse momenta we can write the following expression for polarization $P_\Lambda(s,\xi)$:

$$P_\Lambda(s,\xi) \propto -x_F r^{\text{flip}}_{U\to S} \frac{M}{\zeta} \frac{\int\limits_{b>R(s)} b\,db\, I_0(s,b,\xi)\tilde{U}(s,b)}{\int\limits_{b>R(s)} b\,db\, I_0(s,b,\xi)}, \tag{21}$$

where $R(s) \propto \ln s$ is the hadron interaction radius, which serves as a scale of large and small impact parameter values. Polarization dependence in this region is determined by the unknown function $I_0(s,b,\xi)$ and can have significant values at the transverse momentum which correspond to scattering in the boundary region of the effective field.

5. INCLUSIVE CROSS SECTIONS OF THE UNPOLARIZED HADRON PRODUCTION

To demonstrate self-consistency of the model we consider in this section the unpolarized cross section of Λ- and pion-production processes:

$$\frac{d\sigma}{d\xi} = 8\pi \int\limits_0^\infty b\,db\,\frac{I_0(s,b,\xi)}{|1-iU(s,b)|^2}. \tag{22}$$

At the beginning we approach the process of Λ production. In the fragmentation region we can simplify the problem and consider the process of Λ 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

$$F(s,p_T,x_F)=\frac{is}{x_F^2\pi^2}\int\limits_0^\infty b\,db\,J_0\left(\frac{bp_T}{x_F}\right)\frac{I_0^{1/2}(s,b,x_F)}{1+U(s,b)}. \qquad (23)$$

To obtain Eq. (23) we have used relations $b\simeq x_F\tilde{b}$ and due to the fact that the function I_0 is quadratic on matrix U (the multiparticle analog of the generalazed reaction) the relation follows:

$$I_0^{1/2}(s,b,p_T,x_F)=\frac{s}{\pi^2}\int\limits_0^\infty I_0^{1/2}(s,b,\tilde{b},x_F)J_0(\tilde{b}p_T)\tilde{b}\,d\tilde{b}. \qquad (24)$$

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

$$I_0^{U\to S}(s,b,x_F)=\left[\prod_{Q=1}^{N-1}\langle f_Q(s,b)\rangle\right]\langle f_{U\to S}(s,b,x_F)\rangle, \qquad (25)$$

where N is the total number of quarks in the colliding hadrons.

In the model, the b dependences of the amplitudes $\langle f_Q\rangle$ and $\langle f_{U\to S}\rangle$ 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 by factor $\kappa<1$ (which takes into account reduction of the radius due to spin flip), $r_{U\to S}^{\text{flip}}=\kappa\zeta/\tilde{m}_Q$, and the corresponding function $f_{U\to S}(s,\,b,\,x_F)$ has the form

$$f_{U\to S}(s,b,x_F)=g_{\text{flip}}(x_F)\exp\left(-\frac{\tilde{m}_Q}{\kappa\zeta}b\right). \qquad (26)$$

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

$$I_0(s,b,x_F)=\frac{\bar{g}(x_F)}{g_Q(s)}U(s,b)\exp\left(-\frac{\Delta m_Q b}{\zeta}\right), \qquad (27)$$

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_T and x_F dependences of the Λ-hyperon production cross section and we start with angular distribution*. The corresponding amplitude

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

$F(s,\,p_T,\,x_F)$ 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 β plane and transform the Fourier–Bessel integral over impact parameter into the integral in the complex β plane over the contour C which goes around the positive semiaxis. The amplitude $F(s,\,\beta,\,x_F)$ has the poles and a branching point (at $\beta = 0$) and therefore the amplitude $F(s,\,p_T,\,x_F)$ can be represented as a sum of the poles contribution and the contribution of the cut:

$$F(s,\,p_T,\,x_F) = F_p(s,\,p_T,\,x_F) + F_c(s,\,p_T,\,x_F). \tag{28}$$

Calculation of poles and cut contributions is similar to the case of elastic scattering [37].

The poles and cut contributions determine the inclusive cross-section behaviour of Λ production at moderate and large values of p_T correspondingly, i.e., it will have in the region of large p_T power-like dependence on p_T:

$$\frac{d\sigma}{d\xi} \propto G_c^2(s,x_F)\left(1+\frac{p_T^2}{x_F^2\bar{M}^2}\right)^{-3}, \tag{29}$$

while at smaller values of p_T inclusive cross section would have the exponential p_T dependence:

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

The data for the Λ hyperon production are available at the moderate values of p_T, and the experimental fits to the data [38] of the form

$$A(1-x_F)^n \mathrm{e}^{-B(x_F)p_T}$$

just follow Eq. (30) when relevant parameterization for the function $\bar{g}(x_F)$ is chosen. At high values of p_T, power-like dependence should take place according to Eq. (29). In the energy region of $\sqrt{s} \leqslant 2$ TeV the functions G_p and G_c have very slow variation with energy due to the numerical values of parameters [39].

We can treat the inclusive cross section of the pion production processes in a similar way. In the fragmentation region at small p_T the poles in impact parameter plane at $b \sim R(s)$ lead to the exponential p_T dependence of inclusive cross section. At high p_T the power-like dependence p_T^{-n} with $n = 6$ should take place. It originates from the singularity at zero impact parameter $b = 0$. The exponent n does not depend on x_F. The data are in good agreement with the p_T^{-6} dependence of the unpolarized inclusive cross section (Fig. 7). Recently a similar p_T^{-6} dependence has been obtained for the soft contribution to quark–quark scattering induced by an anomalous chromomagnetic interaction due to instanton mechanism [40].

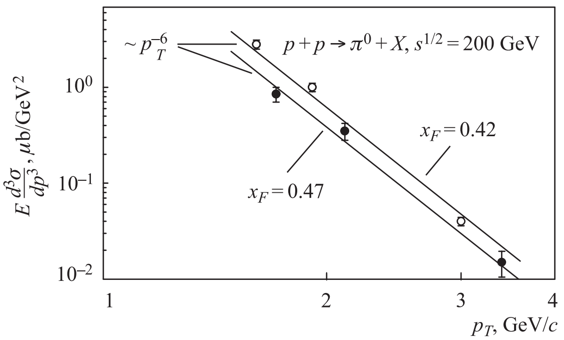

Fig. 7. Transverse momentum dependence of unpolarized inclusive cross section, experimental data from [11]

Thus, in the approach with effective degrees of freedom — constituent quarks and Goldstone bosons — differential cross section at high transverse momenta has a generic power-like dependences on p_T domain.

CONCLUSION

The considered mechanism of SSA generation deals with the effective degrees of freedom and takes into account collective aspects of QCD dynamics. Combined with unitarity, which is an essential part of the approach, it allows one to get a qualitative explanation of the observed regularities: linear dependence on x_F and flat dependence on transverse momentum at large p_T of SSAs in the polarized proton fragmentation region. The spin filtering is responsible for the generation of hyperon polarization in the collisions of the unpolarized nucleons. In particular, it leads to the similar behaviour of Λ polarization. The application of spin filtering to other hyperons is a more complicated case, since those hyperons have two or three strange quarks and the spins of U and D quarks also make contributions into their polarizations.

We also discussed here particle production in the fragmentation region and have shown that the power-like behavior of the differential cross sections at large transverse momenta can be obtained in the approach which has a nonperturbative origin. It is no need to comment that such a dependence always being considered as a manifestation of the genuine hard, short distance processes where asymptotic freedom is at work. Power-like behavior of inclusive cross sections and the strongly interacting nature of quark matter revealed at RHIC, in principle, can be attributed to a different dynamics. However, it is difficult to imagine how both the phenomena can coexist in the strongly interacting coherent medium observed at RHIC when thermalization occurs at very early stage of reaction. It seems natural to suppose that they should have the same origin. One should arrive then to conclusion that the power-like dependence of the differential cross sections should not necessarily be associated with the processes treated by perturbative QCD. This viewpoint gets support in the results on polarization measurements which also indicate possibility of power-like behavior due to soft dynamics. It should also be recollected that the energies where power-like dependence in exclusive processes was observed are evidently too low to be considered as a true asymptotic perturbative QCD. This regime should occur at much higher values of the transverse momenta and energy.

Finally, one should note that in the central and backward regions where correlations between impact parameters of the initial and final particles are weak or even completely degraded, the asymmetry cannot be generated due to the considered mechanism. The experimentally observed vanishing asymmetries in the central and backward regions provide indirect evidence in their favor.

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