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DIFFERENT BEHAVIOUR OF THE SPIN STRUCTURE FUNCTIONS $g_1(x)$ AND $h_1(x)$ AT $x \to 0$

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Abstract

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We consider a low-x behaviour of the spin structure functions $g_1(x)$ and $h_1(x)$ in the unitarized chiral quark model which combines ideas on the constituent quark structure of hadrons with a geometrical scattering picture and unitarity. A nondiffractive singular low-x dependence of $g_1^p(x)$ and $g_1^n(x)$ indicated by the recent SMC experimental data is described. A diffractive type smooth behaviour of $h_1(x)$ is predicted at small x. The expectations for the double-spin asymmetries in the low-mass Drell-Yan production at RHIC in the central region are discussed alongside. PACS: 11.80.Fv, 13.60.Hb, 13.88.+e spin structure functions, unitarity, low-x.

Аннотация

Трошин С.М. и Тюрин Н.Е. Разница в поведении спиновых структурных функций $g_1(x)$ и $h_1(x)$ при $x \to 0$: Препринт ИФВЭ 97-31. – Протвино, 1997. – 10 с., 2 рис., библиогр.: 20.

Рассмотрено поведение структурных функций $g_1(x)$ и $h_1(x)$ в рамках унитаризованной киральной кварковой модели. Описано наблюдаемое коллаборацией SMC недифракционное сингулярное поведение структурных функций $g_1^P(x)$ и $g_1^n(x)$ при малых x. Предсказывается гладкое поведение $h_1(x)$, имеющее при малых x дифракционное происхождение. Обращаются также двойные спиновые асимметрии в процессе Дрелла-Яна с малыми массами.

Introduction

A low-x behaviour of the spin structure functions $g_1(x)$ and $h_1(x)$ is important for understanding a nucleon spin structure. Experimental evaluation of the first moment of g_1 (and the total hadron helicity carried by quarks) is sensitive to extrapolation of $g_1(x)$ to $x = 0$. This extrapolation is a nontrivial issue in the light of the recent CERN SMC [1] and SLAC E154 [2] data indicating a rising behaviour of $g_1(x)$ at small x.

Usually a smooth Regge extrapolation (corresponding to the contribution of the Regge pole of unnatural parity) $g_1 \sim x^{-\alpha_{a_1}}$ $(-0.5 < \alpha_{a_1} < 0)$ being used for extrapolation of the data to $x = 0$. It is evident that the CERN SMC and SLAC E154 experiments give a little support for such smooth behaviour [3,4] and suggest instead a singular dependence on x at $x \to 0$.

A non-Regge behaviour of $g_1(x)$ at small x and the connection of such behaviour with a diffractive contribution is under discussion [5] since the first EMC data have been published [6]. There were used various explicit functional dependencies based either on the perturbative QCD considerations [3,7] or the nonperturbative approaches [8,9]. These parametrizations can successfully describe an increasing behaviour of $g_1^p(x)$ at small values of x. The upper bounds for the low-x behaviour of the spin structure functions were also obtained [8,10] using the general properties of the scattering matrix and some heuristic considerations.

However, the above approaches relate this increasing behaviour with the diffractive contribution to $g_1(x)$ at small x, which does not seem to find the confirmation in the recent experiments. In particular, the SMC data demonstrate the following equality[1]in the region of $0.003 \leq x \leq 0.1$:

$$g_1^p(x) = -g_1^n(x). \tag{1}$$

It might happen that the leading diffractive contribution providing the same sign input to the proton and neutron structure functions would not be able to serve as a sole explanation

[1]Preliminary SMC results [11] based on the data obtained in 1996 indicate a possibility for a more complicated behaviour of the spin structure function $g_1^p(x)$.

of the observed experimental regularities at $x \to 0$, in particular, Eq. (1). Combining the existing facts one could assume the importance of a nondiffractive behaviour of $g_1(x)$ at small x.

The essential point is that the space-time structure of the scattering at small values of x involves the large distances $l \sim 1/Mx$ on the light–cone [12] and the region $x \sim 0$ is therefore sensitive to the nonperturbative dynamics. It leads to the necessity of applying the nonperturbative model approaches.

One should note that the general principles such as unitarity and analyticity are useful and provide some constraints. In particular, unitarity provides the following upper bounds for $g_1(x)$ and $h_1(x)$ [10]:

$$g_1(x) \leq \frac{1}{x}\ln(1/x) \quad \text{and} \quad h_1(x) \leq \frac{1}{x}\ln(1/x). \tag{2}$$

The experimental data for $\Delta\sigma_L(s)$ and $\Delta\sigma_T(s)$ could also be a useful source of information on the low-x behaviour of the spin structure functions. Unfortunately, only the low energy data are available at the moment [13]. Were the experimentally observed decreasing behaviour of $\Delta\sigma_L(s)$ and $\Delta\sigma_T(s)$ also valid at high energies, it would be possible to conclude that at $x \to 0$

$$xg_1(x) \to 0 \quad \text{and} \quad xh_1(x) \to 0. \tag{3}$$

In this paper we show that the non-Regge, nondiffractive behaviour of $g_1(x)$ can be described in the unitarized chiral quark model [14], which combines the ideas on the constituent quark structure of hadrons with a geometrical scattering picture and unitarity. Different functional dependence is predicted for the structure function $h_1(x)$ at small x, which contrary to g_1 has a diffractive origin and satisfies the equality $h_1^p(x) = h_1^n(x)$ at small x. For that purpose we consider the corresponding quark spin densities $\Delta q(x)$ and $\delta q(x)$. We discuss also some predictions which might be interesting for the forthcoming RHIC spin experiments. Namely, the quark spin densities $\Delta q(x)$ and $\delta q(x)$ at small x determine the behaviour of spin asymmetries in hadron-hadron interactions, in particular, double–spin asymmetries A_{LL} and A_{TT} which are to be measured at RHIC in Drell-Yan processes with low-mass lepton pairs.

1. Outline of the model

To obtain the explicit forms for the quark spin densities $\Delta q(x)$ and $\delta q(x)$ it is convenient to use the relations between these functions and discontinuities in the helicity amplitudes of the forward antiquark–hadron scattering [15], which are based on the dominance of the "handbag" diagrams in deep–inelastic processes

$$\begin{aligned} q(x) &= \frac{1}{2}\mathrm{Im}[F_1(s,t) + F_3(s,t)]|_{t=0}, \\ \Delta q(x) &= \frac{1}{2}\mathrm{Im}[F_3(s,t) - F_1(s,t)]|_{t=0}, \\ \delta q(x) &= \frac{1}{2}\mathrm{Im}F_2(s,t)|_{t=0}, \end{aligned} \tag{4}$$

where $s \simeq Q^2/x$ and F_i are the helicity amplitudes for the elastic quark–hadron scattering in the notations for the nucleon–nucleon scattering, i.e.

$$F_1 \equiv F_{1/2,1/2,1/2,1/2},\ F_2 \equiv F_{1/2,1/2,-1/2,-1/2},\ F_3 \equiv F_{1/2,-1/2,1/2,-1/2},\ F_4 \equiv F_{1/2,-1/2,-1/2,1/2}$$

and

$$F_5 \equiv F_{1/2,1/2,1/2,-1/2}.$$

We consider a quark as a structured hadronlike object, since at small x the photon converts to a quark pair at large distances before it interacts with the hadron. At large distances the perturbative QCD vacuum undergoes transition into a nonperturbative one with formation of the quark condensate. The appearance of the condensate means the spontaneous chiral symmetry breaking and the current quark transforms into a massive quasiparticle state – a constituent quark. The constituent quark is embedded into the nonperturbative vacuum (condensate) and therefore we treat it similar to a hadron. The arguments in favour of such a picture can be found in [12,14,16]. The quark–hadron scattering at small x can be considered similar to the hadron–hadron scattering.

The unitary representation for the helicity amplitudes follows from their relations [17] to the U–matrix. In the impact parameter representation:

$$F_{\Lambda_1,\lambda_1,\Lambda_2,\lambda_2}(s,b) = U_{\Lambda_1,\lambda_1,\Lambda_2,\lambda_2}(s,b) + i\rho(s)\sum_{\mu,\nu} U_{\Lambda_1,\lambda_1,\mu,\nu}(s,b)F_{\mu,\nu,\Lambda_2,\lambda_2}(s,b), \tag{5}$$

where λ_i and Λ_i are the quark and hadron helicities, respectively, and b is the impact parameter of quark-hadron scattering. The kinematical factor $\rho(s)$ is unity at high energies. The explicit solution of Eqs. (5) has a rather complicated form, however, in the approximation when the helicity-flip functions are less than the helicity nonflip ones, we can get simple expressions

$$F_{1,3}(s,b) = U_{1,3}(s,b)/[1 - iU_{1,3}(s,b)], \tag{6}$$
$$F_2(s,b) = U_2(s,b)/[1 - iU_1(s,b)]^2. \tag{7}$$

The unitarity requires $\mathrm{Im}U_{1,3}(s,b) \geq 0$. The functions $F_i(s,t)$ are the corresponding Fourier–Bessel transforms of the functions $F_i(s,b)$:

$$F_i(s,t) = \frac{s}{\pi^2}\int_0^\infty bdbF_i(s,b)J_0(b\sqrt{-t}), \tag{8}$$

where $i = 1,2,3$.

We consider now the main points of model [14], which allows one to get an explicit form for the U–matrix. A hadron consists of constituent quarks located at the central part embedded into a nonperturbative vacuum (quark condensate). Justification for this picture is the effective QCD approach and, in particular, the Nambu–Jona-Lasinio model with the six–fermion $U(1)_A$–breaking term. The constituent quark masses can be expressed in terms of the quark condensates. Such quark appears as a quasiparticle, i.e. as a current quark and the surrounding cloud of the quark–antiquark pairs of different flavours. Quantum numbers of the respective constituent quarks are the same as the quantum numbers of current quarks due to conservation of the corresponding currents in QCD. The only exception is the flavour–singlet, axial–vector current. Quark radii are determined by the sizes of the respective clouds. We take the strong interaction radius of quark Q as its Compton wavelength: $r_Q = \xi/m_Q$, where constant ξ is universal for different flavours. Quark formfactor $F_Q(q)$ is taken in the dipole form and the corresponding quark matter distribution $d_Q(b)$ has form [14] $d_Q(b) \propto \exp(-m_Q b/\xi)$. A helicity of the constituent quark J_U in this approach is given by the following sum

$$J_U = 1/2 = J_{u_v} + J_{\{\bar{q}q\}} + \langle L_{\{\bar{q}q\}}\rangle = 1/2 + J_{\{\bar{q}q\}} + \langle L_{\{\bar{q}q\}}\rangle. \quad (9)$$

The current quark helicity is reduced by the negative contribution of the helicities of quarks from the cloud in the structure of constituent quark. So, we can say that a significant part of the spin of constituent quark is associated with the orbital angular momentum of quarks inside this constituent quark, i.e. the cloud quarks rotate coherently inside their constituent quark. Thus, we assume the standard $SU(6)$ spin structure of a nucleon consisting of the three constituent quarks (embedded into the condensate), i.e. the whole nucleon spin is composed of the spins of the constituent quarks. The constituent quarks, however, have a complex internal spatial and spin structure.

The picture of hadron consisting of constituent quarks and surrounding condensate implies that the overlapping and interaction of peripheral clouds occur at the first stage of hadron interaction. Condensates in the overlapping region are excited and, as a result, the quasiparticles , i. e. massive quarks appear in the overlapping region. It should be noted that the condensate excitations are massive quarks, since the vacuum is nonperturbative one and there is no overlap between the physical (nonperturbative) and bare (perturbative) vacuum. Part of hadron energy carried by the outer clouds of condensates being released in the overlapping region, goes to the generation of massive quarks. The number of such quarks fluctuates. Their average number in the framework of the geometrical picture can be estimated as follows:

$$N(s,b) \propto N(s) \cdot D_c^{h_1} \otimes D_c^{h_2}. \quad (10)$$

The function $D_c^{h_i}$ describes the condensate distribution inside hadron h_i; b is the impact parameter of colliding hadrons h_1 and h_2. For the function $N(s)$ we use the maximal possible value $N(s) \propto \sqrt{s}$. Thus, $N(s,b)$ of massive virtual quarks appear in the overlapping region and they generate an effective field. There is another part of the effective field associated with the self-consistent field of the constituent quarks.

Constituent quarks are supposed to scatter in a quasi-independent way by this effective two-component field. We can write down an explicit form for the function $U(s,b)$ for quark-hadron scattering. For simplicity we consider a helicity nonflip part of U–matrix as a pure imaginary one. In accordance with the quasi-independence of valence quarks we represent the input dynamical quantity in the form of product [14]

$$U(s,b) = \prod_{Q=1}^{n+1} f_Q(s,b) \quad (11)$$

in the impact parameter representation. The factors $f_Q(s,b)$ in Eq. 11 correspond to the individual quark amplitudes. The function $f_Q(s,b)$ describes the scattering of a single valence quark in the effective field generated by virtual and valence quarks as it has been discussed above. Combining the above arguments about constituent quark scattering the function $U(s,b)$ can be written as follows [14]:

$$U(s,b) = i\left[\gamma n + \frac{a\sqrt{s}}{\langle m_Q\rangle}\right]^{n+1} \exp(-Mb/\xi). \tag{12}$$

In Eq. (12) n is the number of constituent quarks in the hadron, $M = \sum_{Q=1}^{n+1} m_Q$, and b is the impact parameter of the colliding quark and hadron [14]. a is a universal constant for different hadrons while the constant γ depends on the types of the colliding hadrons.

A mechanism of quark helicity flip in this picture is associated with the constituent quark interaction with the quark generated under interaction of condensates [17]. The quark exchange process between the valence quark and an appropriate quark with relevant orientation of its spin and the same flavour will provide the necessary helicity flip transition, i.e. $Q_+ \to Q_-$.

Of course, such processes are relatively suppressed in comparison with the quark scattering preserving helicity and, therefore, do not contribute into to the spin-averaged observables at the leading order. However, the measuring of the transverse spin asymmetries is sensitive to the subleading contribution and serves, therefore, as a filter of the quark exchange processes $Q_+ \to Q_-$. This transition occurs, when the valence quark knocks out a quark with the opposite helicity and the same flavour. This interaction has the energy suppression by the factor of $1/N(s)$ and it results in the following relation between the functions $U_1(s,b)$ and $U_2(s,b)$:

$$\frac{U_2(s,b)}{U_1(s,b)} \sim \frac{\langle m_Q\rangle^2}{s} \exp[2(\alpha-1)\langle m_Q\rangle b/\xi]. \tag{13}$$

The value of the parameter $\alpha > 1$ reflects the more central character of the quark helicity flip mechanism.

Now we return to the considerations of the helicity nonflip functions U_1 and U_3. These functions differ in the helicities of the initial and final states, but both describe helicity nonflip scatterings. The second term in the square brackets of Eq. (12) results from the quark interaction with the component of the effective field which in its turn arises from the interaction of the condensates. Since the hadron spin is composed of the spins of the constituent quarks, this part of interaction does not depend on the quark spin orientation and the second term in the square bracket in Eq. (12) should be the same for the functions U_1 and U_3 due to the parity conservation. This argument does not work for the first term, which follows from the quark interaction with self-consistent field of the valence quarks. Thus, we should suppose that this part of the effective field depends on the relative spin orientations of the valence quarks and consequently, the constant γ, in fact, has different values in the expressions for the functions U_1 and U_3. It is also flavour dependent.

Thus, instead of Eq. (12) we have

$$U_{1,3}(s,b) = i\left[\gamma_{1,3}n + \frac{a\sqrt{s}}{\langle m_Q\rangle}\right]^{n+1}\exp(-Mb/\xi). \tag{14}$$

Then Eq. (4) allows one to obtain the quark densities $\Delta q(x)$ and $\delta q(x)$ at small x.

2. Results and discussion

Calculating $F_{1,2,3}(s,t)|_{t=0}$ at high values of s, we get $q(x)$, $\Delta q(x)$ and $\delta q(x)$ at small $x \simeq Q^2/s$, i.e.

$$q(x) \sim \frac{1}{x}\ln^2(1/x), \tag{15}$$

$$\Delta q(x) \sim \frac{1}{\sqrt{x}}\ln(1/x) \tag{16}$$

and

$$\delta q(x) \sim x^{\frac{\alpha-1}{n+1}}\ln(1/x). \tag{17}$$

The behaviour of $q(x)$ (and $F_1(x)$), $\delta q(x)$ (and $h_1(x)$) is determined by the leading terms in $U_{1,3}(s,b)$ and $U_2(s,b)$ and therefore the small-x dependence of $F_1(x)$ and $h_1(x)$ will be the universal for the proton and neutron, i.e.:

$$F_1^p(x) = F_1^n(x) \sim \frac{1}{x}\ln^2(1/x). \tag{18}$$

and

$$h_1^p(x) = h_1^n(x) \sim x^{\frac{\alpha-1}{n+1}}\ln(1/x). \tag{19}$$

These functions result from the leading terms and have the features of diffractive origin. It is also seen that $h_1(x)$ has a smooth behaviour at $x \to 0$, i.e. $h_1(x) \to 0$ in this limit ($\alpha > 1$).

The behaviour of $\Delta q(x)$ and correspondingly $g_1(x)$ is determined by the difference $U_3(s,b) - U_1(s,b)$ or the subleading terms in $U_{1,3}$ and therefore the small-x dependence of $g_1(x)$ in this model shows different constant factors for the proton and neutron, i.e.,

$$g_1^{p,n}(x) \simeq \frac{C^{p,n}}{\sqrt{x}}\ln(1/x). \tag{20}$$

Contrary to h_1 the spin structure function g_1 has a singular behaviour at $x \to 0$.

The magnitude and sign of the constants $C^{p,n}$ do not follow from the model. Fit to the SMC data provides a good agreement of Eqs. (20) with experiment at small x ($0 < x < 0.1$) (cf. Fig. 1,2) and also leads to $C^p = -C^n = 0.021$.

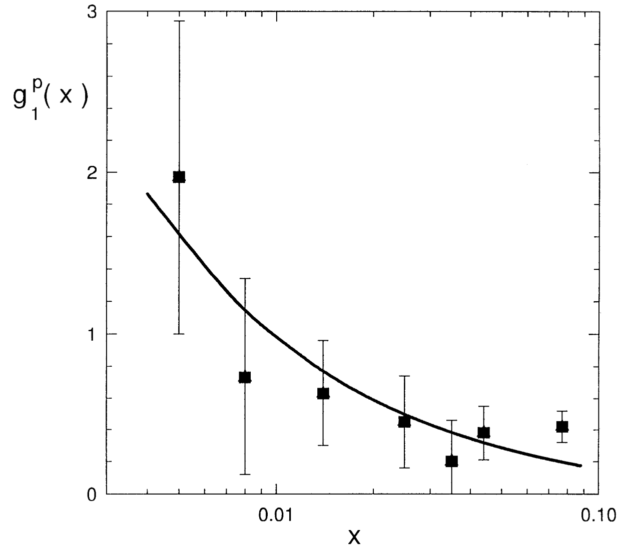

Fig. 1. The x–dependence of the proton spin structure function $g_1^p(x)$.

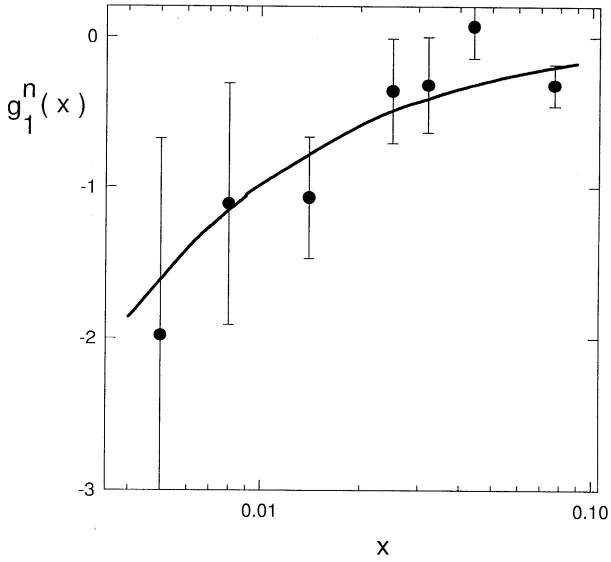

Fig. 2. The x–dependence of the neutron spin structure function $g_1^n(x)$.

The first moment

$$\Gamma_1 = \int_0^1 g_1(x)dx \tag{21}$$

exists. To evaluate Γ_1^p, we use the value of

$$I(0.1, 1) = \int_{0.1}^1 g_1^p(x)dx = 0.092, \tag{22}$$

obtained with the standard parameterization of the data at medium and large values of x [8]. Eq. (20) then gives

$$I(0, 0.1) = \int_0^{0.1} g_1^p(x)dx = 0.057 \tag{23}$$

and for the first moment Γ_1^p, we obtain

$$\Gamma_1^p = \int_0^1 g_1^p(x)dx = 0.149. \tag{24}$$

The above simple estimation of Γ_1^p alongside with the known values of $g_A = 1.257$ and $3F - D = 0.579$ provide the following approximate values for the quark spin contributions:

$$\Delta\Sigma \simeq 0.25, \quad \Delta u \simeq 0.81, \quad \Delta d \simeq -0.45, \quad \Delta s \simeq -0.11, \tag{25}$$

which are in agreement with the results obtained in the comprehensive analysis with account for the QCD evolution and higher twist contributions [1,3,18]. Eq. (25) demonstrates that the singular behaviour of $g_1^p(x)$ in the form of Eq. (20) does not lead to significant deviations from the results of the experimental analysis [1], where the smooth extrapolation of the data to $x = 0$ is used. Nevertheless, the available SMC data in the region of $x \sim 10^{-3}$ provide a reasonable constraint for the functional form of the spin structure function g_1 at small x. We would like to note that the functional dependence of $g_1(x) \sim \frac{1}{\sqrt{x}}\ln(1/x)$ is in agreement with the E154 data as well.

It is important to note that Eqs. (16) and (17) demonstrate that the usual assumption $\Delta q(x) \simeq \delta q(x)$ is not valid at small x. The explicit forms of Eqs. (16) and (17) lead to the conclusion that the longitudinal and transverse double-spin asymmetries for the Drell-Yan processes at $x_F \simeq 0$ are small, i.e.,

$$A_{LL}^{l\bar{l}} \simeq 0 \quad \text{and} \quad A_{TT}^{l\bar{l}} \simeq 0, \tag{26}$$

when the invariant mass of the lepton pair $M_{l\bar{l}}^2 \ll s$. It is just this kinematical region, which is sensitive to low-x behaviour of the spin densities $\Delta q(x)$ and $\delta q(x)$ since:

$$A_{LL}^{l\bar{l}} = -\frac{\sum_i e_i^2[\Delta q_i(x_1)\Delta\bar{q}(x_2) + \Delta\bar{q}_i(x_1)\Delta q_i(x_2)]}{\sum_i e_i^2[q_i(x_1)\bar{q}_i(x_2) + \bar{q}_i(x_1)q_i(x_2)]} \tag{27}$$

and

$$A_{TT}^{l\bar{l}} = a_{TT}\frac{\sum_i e_i^2[\delta q_i(x_1)\delta\bar{q}(x_2) + \delta\bar{q}_i(x_1)\delta q_i(x_2)]}{\sum_i e_i^2[q_i(x_1)\bar{q}_i(x_2) + \bar{q}_i(x_1)q_i(x_2)]}, \tag{28}$$

where $x_1x_2 = M_{l\bar{l}}^2/s$ and $x_F = x_1 - x_2$.

The ratio of the asymmetries $A_{TT}^{l\bar{l}}$ and $A_{LL}^{l\bar{l}}$ is also small in the central region of low-mass Drell-Yan production:

$$A_{TT}^{l\bar{l}}/A_{LL}^{l\bar{l}} \simeq 0. \tag{29}$$

This result agrees with the predictions made in [19].

The above results were obtained in the limit $s \to \infty$ which corresponds to the limit $x \to 0$, i.e. they have an asymptotic nature. However, it might happen that the kinematical region of the SMC experiment lies in the preasymptotic domain and the above formulas in fact are valid at much smaller values of x, than the range covered by the present experiments.

Indeed, it was shown that the preasymptotic effects are very important to understand the experimental regularities observed in the unpolarized scattering [20].

The direct calculation of the quark densities at small x in the preasymptotic region $x > x_0$ according to Eqs. (4,7) and (14) provides the following expressions:

$$q(x) \simeq \frac{c_1}{x}(1 + \frac{c_2}{\sqrt{x}}), \tag{30}$$

$$\Delta q(x) \simeq \frac{c_3}{x}(1 + \frac{c_4}{\sqrt{x}}) \tag{31}$$

and

$$\delta q(x) \simeq c_5(1 + \frac{c_6}{\sqrt{x}}), \tag{32}$$

where the parameters $c_{1,2} \geq 0$. The parameters c_i, $(i = 3 - 6)$ can get negative as well as positive values. It means that in the preasymptotic region the structure functions

g_1 and h_1 may demonstrate a sign changing behaviour. Using the above expressions we arrive again at small values of transverse double-spin asymmetries $A_{TT}^{l\bar{l}}$ in the central region of low-mass Drell-Yan production, but the corresponding longitudinal asymmetries would not disappear in the preasymptotic region. Thus, the measurements of the double-spin asymmetries at RHIC would provide the important information on the x-dependence of spin quark densities $\Delta q(x)$ and $\delta q(x)$.

Conclusion

We considered the low-x behaviour of the spin structure functions $g_1(x)$ and $h_1(x)$ in the framework of the nonperturbative approach on partonic structure of the constituent quarks and their scattering picture in the effective two-component field. We obtained the spin densities of current quarks. Such considerations could be regarded as a kind of a bootstrap approach.

The obtained singular dependence of $g_1(x)$ at $x \to 0$ can describe the rise of the spin structure function observed in the SMC experiment. The model predicts a smooth behaviour of $h_1(x)$ at $x \to 0$ which, contrary to g_1, corresponds to the leading diffractive contribution.

It has been shown that the singular behaviour of the structure function $g_1(x)$ does not affect much the numerical results on the partition of a nucleon spin.

Vanishingly small double-spin asymmetries in the Drell-Yan production with low invariant masses were predicted in the central region.

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