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PARITY VIOLATION IN HIGH-ENERGY PROTON-NUCLEUS SCATTERING AND A QUARK-PARTON MODEL

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Parity violation in high-energy proton-nucleus scattering is calculated using a parton model in which the proton structure is described by Q^2-dependent parton distribution functions. In this model, weak proton scattering is due to quark-quark scattering. Experimental values for scattering of a polarized proton on water can be reproduced with suitable model parameters.

1. Introduction

Convenient theoretical approaches based on meson exchange models[1-3] cannot explain the large parity-violating (PV) effect observed in high-energy proton-nucleus scattering[4]. This prompted additional theoretical efforts. One group of these efforts was aimed at investigating the existence of the parity-non-conserving (pnc) component in the nucleus[5,6]. Another approach[7] used a parton model in order to describe the structure of the nucleon. In that picture, weak nucleon-nucleon scattering was shown to arise from weak quark-quark scattering. Such a picture seems eminently justifiable at high energies. This method has already been extensively studied in a closely related problem of high-transverse-momentum inclusive

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production in polarized proton-proton collisions[8,9]. We apply it here to the elastic scattering of longitudinally polarized protons on nuclei. Our approach is, in principle, similar to the so-called scattering contribution of Ref. 7. However, there are many differences in details and in parametrization, which will be explained in the text.

The scattering contribution, be it described through meson exchange diagrams[1-3] or through weak quark-quark scattering, corresponds to the same basic physics as described by an effective PV weak current-current hamiltonian. Currents are made out of quarks, as described in the general electroweak theory. The difference between a meson exchange model and a quark-parton model is due to the difference in the description and parametrization of the influence of quantum--chromodynamic (QCD) interactions. As QCD effects are not exactly calculable, one has to use some semiempirical model. It is important to find out whether a particular prediction depends on the model used. Moreover, it is possible that one model is better suited for lower energies, while another model could provide a better description of high-energy processes. An additional question is how much some changes in details and in parametrization of a particular type of model can influence the conclusion. Finally, one is also concerned with the predictive power of the theory: is it possible that the nature of the model inevitably leads to too large a number of arbitrary parameters (»hidden« or disclosed)?

The aim of this paper is to give some answers to the questions mentioned above. Our results provide a comparison between two different types of models and also between two different formulations of the quark-parton model for weak scattering. The answers are not irrelevant. If the results of high-energy experiments cannot be ascribed to weak scattering contributions[1-3,7], the only theoretical alternative[5] proposed so far requires careful study[6,10] of the pnc component in the nucleon.

In Sect. 2 we summarize details of our quark-parton model, relying to a great extent on the formalism of Refs. 8 and 9. We also point out how the present theoretical model differs from that used in Ref. 7. In Sect. 3 we display numerical values and discuss the answers to the questions raised. These answers, however, are not quite definite. Empirical uncertainties in the quark-parton model leave some possibility for fitting PV effects. Two parameters, namely an effective gluon mass λ and an effective quark mass m, were needed to fit both the total cross section σ and the asymmetry A. It is encouraging that approximately equal values of parameters, within a factor of $1.3 \sim 2$, can fit two high-energy experiments, one at $p_{lab} = 6\,\mathrm{GeV/c}$ and the other at $E_{lab} = 0.8\,\mathrm{GeV}$. Effective masses are obviously needed in order to parametrize QCD effects which cannot be calculated exactly. Strong nucleon-nucleon scattering, for example, is described by an effective one--gluon exchange[7-9]. This very crude description of the actual physical process has to contain free parameters, i. e. effective masses. As the description of weak parity-violating scattering is also equally naive, fitting parameters are unavoidable.

2. Cross sections in the parton model

Our formalism closely follows that of Refs. 8 and 9, in particular the one exposed in Ref. 9. The only difference is that we deal with the total cross section for scattering of either a polarized or an unpolarized proton on another proton or a neutron.

Scattering between nucleons is understood as scattering between their constituent partons: valence quarks, gluons, quark-antiquark pairs, etc. Parity-violating effects arise from the coherent interference between the weak and strong amplitudes for an individual scattering process such as $ud \to ud$.

We calculate the strong interaction amplitude in an effective gluon model taking into account all possible (anti)quark-(anti)quark scatterings. (Quarks labelled by $u, d, \ldots$ and antiquarks labelled by $\bar{u}, \bar{d}, \ldots$ are understood under the generic name partons.) The elementary partons have momenta

$$p_i + p_j \to p_k + p_l. \tag{2.1}$$

These momenta are connected with the momenta of the initial baryons

$$p_a + p_b \to X \tag{2.2}$$

and with the Mandelstam variables $\hat{s}$, $\hat{t}$ and $\hat{u}$ as follows:

$$\begin{aligned} \hat{s} &= (p_i + p_j)^2 = (x_i p_a + x_j p_b)^2, \\ t &= (p_i - p_k)^2 = (x_i p_a - p_k)^2, \\ \hat{u} &= (p_i - p_l)^2 = (x_i p_a - p_l)^2. \end{aligned} \tag{2.3}$$

A complete list of parton cros-sections and their relations is the following:

$$\begin{aligned} \sigma(uu) &= \sigma_t(uu) + \sigma_u(uu) + \sigma_{tu}(uu), \\ \sigma(dd) &= \sigma_t(dd) + \sigma_u(dd) + \sigma_{tu}(dd), \\ \sigma(u\bar{u}) &= \sigma_t(u\bar{u}) + \sigma_s(u\bar{u}) + \sigma_{ts}(u\bar{u}), \\ \sigma(d\bar{d}) &= \sigma_t(d\bar{d}) + \sigma_s(d\bar{d}) + \sigma_{ts}(d\bar{d}), \\ \sigma(u\bar{d}) &= \sigma_t(u\bar{d}), \\ \sigma(d\bar{u}) &= \sigma_t(d\bar{u}), \\ \sigma(d\bar{u}) &= \sigma_t(d\bar{u}_{,}, \\ \sigma(\overline{uu}) &= \sigma(uu), \\ \sigma\overline{(dd)} &= \sigma(dd), \\ \sigma(ud) &= \sigma_t(ud), \end{aligned}$$

and

$$\begin{aligned} \sigma_t(uu) &= \sigma_t(dd), \\ \sigma_u(uu) &= \sigma_u(dd), \\ \sigma_{tu}(uu) &= \sigma_{tu}(dd). \end{aligned} \tag{2.4}$$

An effective gluon, with an effective mass λ, can be exchanged in either the $\hat{t}$, $\hat{s}$ or $\hat{u}$ channel. Quarks and antiquarks are also given an effective mass m. These effective masses are arbitrary semiempirical parameters. As discussed below, they will be fixed in such a way as to reproduce measured strong cross sections for nucleon-nucleon scattering. As an example we list the total cross sections for the u-flavoured quark-quark scattering:

$$\sigma_t(uu) = \left[\frac{2}{9}\right] \frac{2\pi\alpha_s^2}{s} \left\{1 + \frac{2s(s-4m^2)+8m^4}{\lambda^2(s-4m^2+\lambda^2)} + \right.$$

$$\left. + 2\frac{s}{s-4m^2} \ln \frac{\lambda^2}{s-4m^2+\lambda^2}\right\}, \tag{2.5a}$$

$$\sigma_u(uu) = \left[\frac{2}{9}\right] \frac{2\pi\alpha_s^2}{s} \left\{\frac{2s}{\lambda^2} - \frac{s+4m^2-2\lambda^2}{s-4m^2+\lambda^2} + \right.$$

$$\left. + 2\frac{s+\lambda^2}{s-4m^2} \ln \frac{\lambda^2}{s-4m^2+\lambda^2}\right\}, \tag{2.5b}$$

$$\sigma_{ut}(uu) = \left[\frac{2}{27}\right] \frac{2\pi\alpha_s^2}{s} \frac{s^2-8m^2s+12m^4}{s-4m^2} \times \tag{2.5c}$$

$$\times \frac{2}{s+2\lambda^2-4m^2} \ln \frac{\lambda^2}{s-4m^2+\lambda^2},$$

$$\sigma^{tot}(uu) = \sigma_u(uu) + \sigma_t(uu) - 2\sigma_{ut}(uu). \tag{2.5d}$$

The $u-u$ and $u-\bar{u}$ scatterings are related when the gluon is exchanged in the t-channel

$$|\sigma_t(u\bar{u})|^2 = |\sigma_t(uu)|^2. \tag{2.6}$$

However, the gluon can also be exchanged in the s-channel in the case of $u-\bar{u}$ scattering, so that one has, for example,

$$\sigma_s(u\bar{u}) = \left[\frac{2}{9}\right] \frac{4\pi\alpha_s^2}{3s} \frac{(s+2m^2)^2}{(s-\lambda^2)^2}, \tag{2.7a}$$

$$\sigma_{st}(u\bar{u}) = \left[\frac{2}{27}\right] \frac{2\pi\alpha_s^2}{s} \frac{1}{\lambda^2-s} \left\{\lambda^2 + \frac{3}{2}(s-4m^2) + \right.$$

$$\left. + \frac{(s+\lambda^2)^2-4m^4}{s-4m^2} \ln \frac{\lambda^2}{s-4m^2+\lambda^2}\right\}. \tag{2.7b}$$

In the parton-model sense, the total cross section σ for proton-proton scattering is built up of parton cross sections $\overline{\sigma}$:

$$\sigma = \sum_{i,j,k,l} \int \mathrm{d}x_i \, \mathrm{d}x_j \, \overline{\sigma} \, (ij \to kl) \times \\ \times G_{i/p} (x_i, Q^2) \, G_{j/p} (x_j, Q^2). \tag{2.8}$$

Here, $G_{i/p} (x_i, Q^2)$ are parton distribution functions which determine the number of partons of the type i ($i = u, d, \overline{u}, \ldots$) in a proton. Each parton carries momentum

$$p_i = x_i \, p_a. \tag{2.9}$$

The dependence of $G_{i/p}$ on Q^2 has been obseved in deep-inelastic electron-proton scattering $[e(p) + p(p) \to e(p') + x; \; Q = (p - p')^2]$, and is theoretically associated with the asymptotic freedom in QCD[11–13].

The analog of Q^2 in parton-parton scattering, Eq. (2.1), is not directly observable. Various combinations of $\hat{s}$, $\hat{t}$ and $\hat{u}$ have been suggested, such as

$$Q^2 = \frac{2\hat{s} \; \hat{t} \; \hat{u}}{\hat{s}^2 + \hat{t}^2 + \hat{u}^2} \tag{2.10}$$

or

$$Q^2 = (\hat{s} \; \hat{t} \; \hat{u})^{1/3}, \; Q^2 = \hat{t},$$

$$Q^2 = \hat{s}, \; Q^2 = \frac{\hat{s} - \hat{t} - \hat{u}}{3}.$$

We will return to parton distribution functions after considering the mixing of the weak and strong amplitudes which is responsible for parity-violating asymmetry. The total amplitude is

$$\mathcal{M} = \mathcal{M}_s + \mathcal{M}_w, \tag{2.11}$$

where s denotes the strong, effective gluon exchange contribution, and w is concerned with the intermediate vector boson (IVB) exchanges. The cross section is determined by the absolute square of Eq. (2.11):

$$\sum_{\text{spin}} \sum_{\text{colour}} |\mathcal{M}|^2 = \sum_{\text{spin}} \sum_{\text{colour}} (|\mathcal{M}_s|^2 + 2 \, \mathrm{Re} \, \mathcal{M}_s \mathcal{M}_w^* + |\mathcal{M}_w|^2). \tag{2.12}$$

Thus,

$$\hat{\sigma} = \hat{\sigma}_s + \hat{\sigma}_w + \hat{\sigma}_{sw}. \tag{2.13}$$

Here, $\hat{\sigma}_w \sim G_F^2$ can be neglected. In order to find parity-violating asymmetry, one has to calculate σ_{sw}, for which we use the standard electroweak model:

$$H_{ew} = \sum_{i=e,u,d,s,c} [e \, Q_i \, \overline{\psi}_i \, \gamma^\mu \, \psi_i \, A_\mu + \overline{\psi}_i \, \gamma^\mu \, (g_V^i - g_A^i \, \gamma_5) \, \psi_i \, Z_\mu] + \\ + \sum_{\substack{i=e,d,s \\ j=\nu_e,u,c}} G_{V-A}^{ij} \, [\overline{\psi}_i \, \gamma^\mu \, (1 - \gamma_5) \, \psi_j \, W_\mu + H. \, C.], \tag{2.14}$$

where

$$Q_e = -1, \quad Q_u = Q_c = \frac{2}{3}, \quad Q_d = Q_s = -\frac{1}{3}$$

and

$$G^{ij}_{V-A} = \frac{M_W \sqrt{G_F}}{2^{1/4}} [\delta_{le}\, \delta_{j\nu e} + (\cos\Theta_c\, \delta_{id} + \sin\Theta_c \delta_{is})\, \delta_{ju}$$

$$+ (\cos\Theta_c\, \delta_{is} - \sin\Theta_c\, \delta_{id})\, \delta_{jc}],$$

$$\sin^2\Theta_c \approxeq 0.06, \; G_F \approxeq 1.01 \times 10^{-5}\, m_p^{-2} \simeq 10^{-5}\, \mathrm{GeV}^{-2},$$

$$\begin{bmatrix} g^l_V \\ g^l_A \end{bmatrix} = 2^{1/4} M_Z \sqrt{G_F} \begin{bmatrix} a_l \\ b_l \end{bmatrix},$$

$$a_e = -\frac{1}{2} + 2x, \qquad a_u = a_c = \frac{1}{2} - \frac{4}{3} x,$$

$$a_d = a_s = -\frac{1}{2} + \frac{2}{3} x, \; b_e = b_d = b_s = -\frac{1}{2}, \; b_u = b_c = \frac{1}{2},$$

$$x \equiv \sin^2\Theta_W \cong 0.224 \pm 0.019.$$

In keeping with the spirit of the parton model[8,9], we have not introduced QCD short-distance corrections. We assume that QCD corrections are semiempirically included in the proton distribution functions $G(x, Q^2)$, so that the introduction of the renormalization-group log-corrections might lead to double counting.

Our expression for the cross section will also depend on IVB masses[14]:

$$M_{W^\pm} = 81.0 \pm 2.5\ \mathrm{GeV/c^2},$$

$$M_{Z^0} = 91.9 \pm 1.3\ \mathrm{GeV/c^2}, \tag{2.15a}$$

and IVB widths[8,9]:

$$\Gamma_Z = \frac{4\hat{s}}{3\pi} \frac{G_F}{\sqrt{2}} M_Z \left(1 - 2\sin^2\Theta_W + \frac{8}{3}\sin^4\Theta_W\right),$$

$$\Gamma_W = \frac{4\hat{s}}{3\pi} \frac{G_F}{\sqrt{2}} M_W. \tag{2.15b}$$

The cross section for polarized protons (nucleons) will be, via formula of the type (2.8), connected with the cross sections for polarized quarks. Denoting the right- (left-)handed quark by R (L), one finds that

$$\sigma^{tot}(u_{R,L}; \bar{u}) = \left[\frac{4}{9}\right] \frac{\alpha_s M_Z^2 G_F}{2\sqrt{2}} \cdot \frac{1}{s(s-4m^2)} \times$$

$$\times \left[\frac{1}{s-\lambda^2}\left\{(a \mp 1)^2 \left[\left[3m^2 s - 2m^4 - s^2 - \right.\right.\right.\right.$$

$$\left. - 2(s-2m^2) M_Z^2 - M_Z^4 \right] \ln \frac{M_Z^2}{s-4m^2+M_Z^2} -$$

$$\left. - (s-4m^2)\left[\frac{1}{2}(s-4m^2) - M_Z^2 - 2(s-2m^2)\right]\right] -$$

$$\left. - (a^2-1) m^2 \left[(s-4m^2+2M_Z^2) \times \ln \frac{M_Z^2}{s-4m^2+M_Z^2} + 2(s-4m^2)\right]\right\} +$$

$$+ \frac{s-M_Z^2}{(s-M_Z^2)^2 + M_Z^2 \Gamma_Z^2}\left\{(a \mp 1)^2 \left[\left[4m^2 s - 2m^4 - s^2 - (2s-3m^2)\lambda^2 - \lambda^4\right] \times \right.\right.$$

$$\left. \times \ln \frac{\lambda^2}{s-4m^2+\lambda^2} + (s-4m^2)\left[\frac{1}{2}(s-4m^2) - \lambda^2 - (2s-3m^2)\right]\right] -$$

$$\left.\left. - (a^2-1) m^2 \left[(2s-4m^2+\lambda^2) \ln \frac{\lambda^2}{s-4m^2+\lambda^2} + (s-4m^2)\right]\right\}\right], \tag{2.16a}$$

$$\sigma^{tot}(u_{R,L}; u) = \left[\frac{4}{9}\right] \frac{\alpha_s G_F M_Z^2}{2\sqrt{2}s(s-4m^2)} \frac{1}{s+\lambda^2+M_Z^2-4m^2} \times$$

$$\times \left\{(a \mp 1)^2 \left[(s - 5m^2 s + 6m^4)\left(\ln \frac{\lambda^2}{s-4m^2+\lambda^2} + \ln \frac{M_Z^2}{s-4m^2+M_Z^2}\right) + \right.\right.$$

$$\left.\left. + m^2 (s-4m^2+\lambda^2) \ln \frac{\lambda^2}{s-4m^2+\lambda^2} - M_Z^2 \ln \frac{M_Z^2}{s-4m^2+M_Z^2}\right]\right] +$$

$$+ \left[(a^2-1) m^2 \left[-s\left(\ln \frac{\lambda^2}{s-4m^2+\lambda^2} + \ln \frac{M_Z^2}{s-4m^2+M_Z^2}\right) + \right.\right.$$

$$\left.\left.\left. + \left[-(s-4m^2+\lambda^2) \ln \frac{\lambda^2}{s-4m^2+\lambda^2} + M_Z^2 \ln \frac{M_Z^2}{s-4m^2+M_Z^2}\right]\right]\right\}, \tag{2.16b}$$

$$\sigma^{tot}(u_L, d) = \left[\frac{4}{9}\right] \frac{4\alpha_s M_W^2 G_F \cos^2\Theta_C}{\sqrt{2}\, s\,(s-4m^2)(s-4m^2+\lambda^2+M_W^2)} \times$$

$$\times \left\{\left[s(4m^2-s)-2m^4\right]\left(-\ln\frac{M_W^2}{s-4m^2+M_W^2}-\ln\frac{\lambda^2}{s-4m^2+\lambda^2}\right)-\right.$$

$$\left.-m^2\left[(s-4m^2+M_W^2)\ln\frac{M_W^2}{s-4m^2+M_W^2}-\lambda^2\ln\frac{\lambda^2}{s-4m^2+\lambda^2}\right]\right\}, \qquad (2.16c)$$

$$\sigma^{tot}(u_L, \bar{d}) = \left[\frac{4}{9}\right] \frac{4\alpha_s \cos^2\Theta_C (s-M_W^2) M_W^2 G_F}{\sqrt{2}\, s\,(s-4m^2)} \times$$

$$\times \frac{1}{(s-M_W^2)^2+M_W^2\Gamma_W^2}\left\{(s-4m^2)\left(m^2-\frac{3}{2}s-\lambda^2\right)+\right.$$

$$\left.+\left[4m^2 s-2m^4+\lambda^2(3m^2-2s)-\lambda^4\right]\ln\frac{\lambda^2}{s-4m^2+\lambda^2}\right. . \qquad (2.16d)$$

Important equalities are the following:

$$\sigma(u_L d) = \sigma(d_L u);\ \sigma(d_{R,L}\bar{d}) = \sigma(d_{R,L}\bar{u})\,|_{a\to b},$$
$$\sigma(u_L \bar{d}) = \sigma(d_L \bar{u});\ \sigma(d_{R,L} d) = \sigma(u_{R,L} u)\,|_{a\to b},$$
$$\sigma(u_R d) = \sigma(u_R \bar{d}) = \sigma(d_R u) = \sigma(d_R \bar{u}) = 0. \qquad (2.16e)$$

Here, a and b are defined by

$$a = 1 - \frac{8}{3}\sin^2\Theta,$$
$$b = 1 - \frac{4}{3}\sin^2\Theta. \qquad (2.17)$$

In order to calculate the asymmetry

$$A = \frac{\sigma_+ - \sigma_-}{\sigma_+ + \sigma_-}, \qquad (12.18)$$

one has to find cross sections for the scattering of the initial proton of positive (negative) helicity. (The cross section (2.8) is $\sigma = \sigma_+ + \sigma_-$.)

$$\sigma^{\pm} = \sum_{i,j,k,l} \int \mathrm{d}x_i\, \mathrm{d}x_j\, [G_{i^+/p^\pm}(x_i, Q^2)\, \hat{\sigma}(i^+ j \to kl)$$
$$+ G_{i^-/p^\pm}(x_i, Q^2)\, \hat{\sigma}(i^- j \to kl)]\, G_{j/p}(x_j, Q^2). \qquad (2.19)$$

The polarized distribution functions such as $G_{i^{\pm}/p^{+}}$ have been studied experimentally in deep-inelastic scattering. We will rely on the fit of Ref. 8:

$$G_{u^+/p^+}(x, Q^2) = \frac{1}{2}(1 + x^{0.39})\, G_{u/p}(x, Q^2),$$

$$G_{u^-/p^+}(x, Q^2) = \frac{1}{2}(1 - x^{0.39})\, G_{u/p}(x, Q^2),$$

$$G_{d^+/p^+}(x, Q^2) = \frac{1}{2}\left(1 - \frac{1}{3}x^{0.23}\right) G_{d/p}(x, Q^2),$$

$$G_{d^-/p^+}(x, Q^2) = \frac{1}{2}\left(1 + \frac{1}{3}x^{0.23}\right) G_{d/p}(x, Q^2). \qquad (2.20)$$

We use the unpolarized distribution function $G_{i/p}$ of Ref. 12. The integration over x_s' s in (2.8) and (2.19) depends on the form of the distribution function and therefore we show it here in detail:

$$G_{d/p}(x, Q^2) = \frac{1}{B\,[(0.85-0.24\,\bar{s}), (4.35+0.816\,\bar{s})]} \times$$

$$\times\, x^{-0.15-0.24\bar{s}}(1-x)^{3.35+0.816\bar{s}}, \qquad (2.21a)$$

$$G_{u/p}(x, Q^2) = \frac{3}{B\,[(0.70-0.176\,\bar{s}), (3.60+0.8\,\bar{s})]} \times$$

$$\times\, x^{-0.30-0.176\bar{s}}(1-x)^{2.60+0.8\bar{s}} - G_{d/p}(x, Q^2), \qquad (2.21b)$$

$$G_{\bar{u}/p}(x, Q^2) = G_{\bar{d}/p}(x, Q^2) = \frac{1}{6} A_S(\bar{s})\, x^{-1}(1-x)^{\eta_s(\bar{s})}. \qquad (2.21c)$$

Here, B symbolizes the Euler beta function

$$B(x, y) = \frac{\Gamma(x)\,\Gamma(y)}{\Gamma(x+y)}, \qquad (2.22)$$

and

$$A_S(\bar{s}) = 1.21 + 1.4\,\bar{s} + 0.6\,\bar{s}^2. \qquad (2.23)$$

The variable $\bar{s}$ is defined by

$$\bar{s} = \ln\frac{\ln Q^2/\Lambda^2}{\ln Q_0^2/\Lambda^2}, \qquad (2.24)$$

$$\Lambda = 0.3 \text{ GeV}, \quad Q_0^2 = 1.8 \text{ GeV}^2.$$

Useful relations are also the following:

$$G_{i^+/p^+} = G_{i^-/p^-},$$

$$G_{i^+/p^-} = G_{i^-/p^+},$$

$$G_{\bar{u}^\pm/p^+} = G_{\bar{d}^\pm/p^+} = \frac{1}{2} G_{\bar{u}/p}. \tag{2.25}$$

The expressions for distribution functions make sense only for $\bar{s}$ finite, which means that

$$\ln Q^2/\Lambda^2 > 0 \text{ or } Q^2 \geq \Lambda^2. \tag{2.26}$$

In our case, this condition is best implemented by choosing from (2.6)

$$Q^2 = \hat{s} \geq \Lambda^2. \tag{2.27}$$

The condition (2.27) avoids singularity, which would appear in Eqs. (2.5), (2.7) and (2.16) for $\hat{s}=0$. On the basis of Eq. (2.3), $\hat{s}=0$ must appear if one allows that each x_i varies in the region $0 \leq x_i \leq 1$. The condition (2.27) excludes the small region around zero in the x_i, x_j plane over which the integration in (2.8) and (2.19) was performed. The condition (2.27) leads to the finite integrals for the total cross sections.

In Sect. 3 we shall comment on the alternative formulation which was employed by Ref. 7. We shall give numerical results and conclusions in Sect. 4. In order to test the dependence of the results for σ and A on the choice of the parton distribution function, we have carried out the analysis also by using the functions presented in Ref. 13. Although these functions do not depend on Q^2, we have performed the integration over x_i and x_j in (2.8) and (2.19), in accordance with the condition (2.27).

3. Comparison with other calculations

The present approach is rather different from the calculations based on meson exchange models[1–3]. The quark structure in these models emerges openly only in the calculation of weak meson vertices. Such calculations are closely connected with the theory of hyperon non-leptonic decays. The present formalism is strongly related with deep-inelastic electron-proton scattering. At present one cannot give a rigorous theoretical argument in favour of one or the other approach. This would require an ability to solve QCD dynamics explicitly. One feels intuitively[7–9] that a parton model might be more suitable for higher energies.

Although in principle equivalent, the present parton-model approach differs in much detail from Ref. 7. In Ref. 7, nucleons are described by SU (6) quark-model wave functions. Orientation of valence quark spins in nucleon wave functions is used to calculate polarizations. Quark flavours f, as they appear in SU (6) wave functions, are weighted by the unpolarized distribution functions. The symbol $F_f(x)$ in the notation of Ref. 7 corresponds to our $G_{f/N}(x, Q^2)$. The authors of Ref. 7 used the distribution functions from Ref. 13, which do not depend on Q^2. In order to describe polarization-related phenomena, we have used the polarized distribution functions (2.20) from Ref. 9, which do depend on Q^2. In that formalism[8, 9], SU (6) wave functions are not used. The Q^2 dependence of the distribution functions by itself (as shown in Sect. 4) does not make much of numerical difference. However, the formalism of Refs. 8 and 9, contrary to that used by Ref. 7, includes, at least partially, sea contributions besides valence quark contributions. On the other hand, Ref. 7 used QCD-corrected quark-quark scattering amplitudes. QCD corrections were calculated both at the one-loop level and in the leading-logarithm approximation using renormalization-group techniques. We have used the bare electroweak lagrangian to describe weak interactions, assuming that appropriate QCD corrections are included[12] in the $G_{s/N}(x, Q^2)$ distribution functions. All this is best illustrated by comparing formulae (2.11) and (A. 1) of Ref. 7 with our formulae (2.5), (2.7) and (2.16). In the present paper, the expressions contain many contributions from quark-antiquark scattering which are absent in Ref. 7. This paper contains nothing equivalent to the one-loop corrections, i. e. terms containing L_i's $(i = 1, \ldots, 6)$ in (A. 1) of Ref. 7. Only the contributions based on the operator-product expansion (OPE) can be correlated with our results.

The expressions of Ref. 7 are calculated in the high-energy limit. If one assumes $\hat{s} \gg$ in (2.5), one immediately finds that

$$\hat{\sigma} \sim \pi\, \alpha_s^2/\lambda^2. \tag{3.1}$$

This constant (or nearly constant) parton cross section leads immediately to a constant cross section for nucleon-nucleon scattering, as displayed in formula (2.11) of Ref. 7. Calculations presented in the next section have been performed numerically by introducing full formulae (2.5) into expression (2.8).

In order to compare the weak-strong mixing, i. e. (A. 1) of Ref. 7 with Eq. (2.16), one has to use the high-energy limit of the Mandelstam variables (2.3); for example,

$$\begin{aligned} \hat{s} &= x_1^2\, p_a^2 + x_2\, p_b^2 + 2x_1\, x_2\, p_a \cdot p_b \\ &\to x_1\, x_2\, s. \end{aligned} \tag{3.2}$$

Here, s refers to nucleon-nucleon scattering, i. e. according to Eq. (2.2):

$$\begin{aligned} s = (p_a + p_b)^2 &= p_a^2 + p_b^2 + 2p_a\, p_b, \\ 2p_a\, p_b &\to s, \\ p_a^2 = p_b^2 &\to 0. \end{aligned} \tag{3.3}$$

Instead of a total parton cross section, for example (2.16b), one has to use the corresponding partial cross scction for $m = 0$[9]:

$$\frac{d\hat{\sigma}}{dt}(u_{R,L}, u) = -\frac{4}{9}\frac{\alpha_s G_F}{2\sqrt{2}}(a \mp 1)^2 \times$$

$$\times \left\{\frac{1}{\hat{t} - \lambda^2(1 - \hat{u}/M_Z^2)} + \frac{1}{\hat{u} - \lambda^2(1 - \hat{t}/M_Z^2)}\right\}. \tag{3.4}$$

This is related to (2.16b) by

$$\hat{\sigma}(u_{R,L}, u)(m = 0) = \int_{-\hat{s}}^{0} \frac{d\hat{\sigma}}{dt} dt. \tag{3.5}$$

In order to calculate the first term in (A. 1a) of Ref. 7, we use

$$M_Z^2 \to \infty \tag{3.6}$$

and neglect the $(u-\lambda^2)^{-1}$ term in the spirit of the high-energy limiting procedure:

$$\hat{u} = -\hat{s} - \hat{t},$$

$$\hat{s} \gg |\hat{u}| \gg |\hat{t}|. \tag{3.7}$$

Using (3.2) one finds that

$$\sigma(u_{R,L}, u) = +\frac{4}{9}\frac{\alpha_s G_F}{2\sqrt{2}}(a \mp 1)^2 \ln\left(1 + \frac{x_1 x_2 s}{\lambda^2}\right). \tag{3.8}$$

The combination

$$\hat{\sigma}(u_L, u) - \hat{\sigma}(u_R, u) = \frac{8}{9} G_F \alpha_s \left(1 - \frac{8}{3}\sin^2\Theta_w\right) \ln\left(1 + \frac{x_1 x_1 s}{\lambda^2}\right) \tag{3.9}$$

is proportional to the first term in (A. 1a) of Ref. 7. The numerical results presented in the next section have been obtained by using complete expressions (2.16) and formula (2.19).

4. *Numerical results and discussion*

All our numerical results depend on two parameters, namely an effective gluon mass λ and an effective quark (parton) mas m. (We have used $\alpha_s = 1$, as in Ref. 7). The experimental strong total cross section for $p-p$ and $p-n$ scattering is fairly constant:

$$\sigma \cong \sigma_+ + \sigma_- \cong 110\,\text{GeV}^{-2}, \tag{4.1}$$

through the energy region spanned by two existing experiments[4]: (4.2)

$$A = (2.65 \pm 0.96) \times 10^{-6} \text{ at } p_{lab} = 6 \text{ GeV/c}, \tag{4.3}$$

$$A = (6.6 \pm 3.2) \times 10^{-7} \text{ at } E_{lab} = 800 \text{ MeV}.$$

The values of $A(pp)$ and $A(pn)$ were calculated for various values of the parameters m and λ. These parameters were selected so that σ always had the value given by Eq. (4.1). In Tables 1, 2, 3 and 4 we give values for A which are reasonably close to the experimental ones. Tables 1 and 2 correspond to the experimental values given by (4.2). Table 1 is calculated by using the Q^2-dependent parton distribution functions of Ref. 12, while Table 2 has been obtained with the distribution functions of Ref. 13. In Tables 3 and 4 we present the same type of data for the experiment (4.3). No significant difference has been found between the results based on two different types of distribution function. However, one should keep in mind that the usage of the Q^2-dependent distribution functions is more consistent with the cut-off based on the condition (2.27).

It is possible to find pairs of m and λ for which one can closely reproduce experimental results. For the experiment (4.2), for example, Table 1 shows $A(pp) = 2.02 \times 10^{-6}$ for $\lambda = 0.0158$ GeV and $m = 0.866$ GeV. Similar but not equal parameter values are also encountered in the experiment (4.3). In this case one finds, for example, $A(pp) = 3.27 \times 10^{-7}$ for $\lambda = 0.0084$ GeV and $m = 0.656$ GeV. It is surprising that so heavy an effective parton mass should be needed, while the effective gluon mass is more than an order of magnitude smaller than the one ($\lambda = 0.5$ GeV) used by Ref. 7. One could say that, in our approach, m has somehow taken over the role played by λ in Ref. 7.

It is to some extent encouraging that both experiments, performed at quite different energies, can be fitted with similar values of λ and m. However, a satisfactory and really convincing theoretical answer would require that both experiments be fitted with the same parameters.

The value of the total cross section depends strongly on λ, as it was found in Ref. 7. Very much smaller values of λ needed here stem from the differences in formalism which we discussed in the preceding section. The difference between the polarized cross sections, $\sigma_+ - \sigma_-$, depends more on m than on λ. This is illustrated in Table 5, which corresponds to the experiment (4.2) and which shows σ and $\sigma_+ - \sigma_-$ as functions of λ for two fixed values of m.

The λ, m dependence displayed in Table 5 qualifies our preceding observation that both experiments can be fitted with similar values of λ and m. This cannot be due to just an accidental coincidence, but probably displays some internal consistence of the model. However, the model is so crude and its physical content is so simplified that one cannot expect an ideal description of the real world.

On the basis of the present analysis, the so-called scattering contribution can be discarded as an adequate description of the large parity-violating effects observed[4] in high-energy proton-nucleus scattering. As the theoretical description depends on two fitting parameters, the case for the scattering contribution cannot be proved beyond doubt.

TABLE 1.

m^2 (GeV^2)	0.6595	0.7495	0.775	0.7865	0.830	0.833	0.952	0.955	1.055
λ^2 (GeV^2)	0.0005	0.00025	0.0002	0.00015	0.00011	0.00012	0.00009	0.0001	0.00007
$A(pp)$	$7.15 \cdot 10^{-6}$	$2.02 \cdot 10^{-6}$	$1.38 \cdot 10^{-6}$	$1.12 \cdot 10^{-6}$	$6.24 \cdot 10^{-7}$	$6.26 \cdot 10^{-7}$	$1.31 \cdot 10^{-7}$	$1.27 \cdot 10^{-7}$	$2.55 \cdot 10^{-8}$
$A(pn)$	$7.17 \cdot 10^{-6}$	$2.1 \cdot 10^{-6}$	$1.4 \cdot 10^{-6}$	$1.12 \cdot 10^{-6}$	$6.31 \cdot 10^{-7}$	$6.4 \cdot 10^{-7}$	$1.45 \cdot 10^{-7}$	$1.44 \cdot 10^{-7}$	$2.85 \cdot 10^{-8}$

Asymmetry at $p_{lab} = 6$ GeV/c; distribution functions from Ref. 12.

TABLE 2.

m^2 (GeV^2)	0.610	0.710	0.739	0.745	0.787	0.790	0.940	0.942	0.977
λ^2 (GeV^2)	0.0005	0.00025	0.0002	0.00015	0.00011	0.00012	0.0009	0.0001	0.00007
$A(pp)$	$7.00 \cdot 10^{-6}$	$1.96 \cdot 10^{-6}$	$1.21 \cdot 10^{-6}$	$1.25 \cdot 10^{-6}$	$7.13 \cdot 10^{-7}$	$7.19 \cdot 10^{-7}$	$1.63 \cdot 10^{-7}$	$1.54 \cdot 10^{-7}$	$8.17 \cdot 10^{-8}$
$A(pn)$	$7.43 \cdot 10^{-6}$	$2.21 \cdot 10^{-6}$	$1.40 \cdot 10^{-6}$	$1.47 \cdot 10^{-6}$	$8.34 \cdot 10^{-6}$	$8.41 \cdot 10^{-7}$	$1.91 \cdot 10^{-7}$	$1.80 \cdot 10^{-7}$	$9.80 \cdot 10^{-8}$

Asymmetry at $p_{lab} = 6$ GeV/c; distribution functions from Ref. 13.

TABLE 3.

m^2 (GeV^2)	0.247	0.296	0.297	0.300	0.350	0.394	0.430	0.460	0.510
λ^2 (GeV^2)	0.0005	0.0002	0.00025	0.00015	0.0001	0.00009	0.00007	0.00005	0.00003
$A(pp)$	$1.80 \cdot 10^{-4}$	$2.65 \cdot 10^{-5}$	$3.03 \cdot 10^{-5}$	$6.06 \cdot 10^{-6}$	$3.70 \cdot 10^{-6}$	$1.06 \cdot 10^{-6}$	$3.27 \cdot 10^{-7}$	$9.41 \cdot 10^{-8}$	$1.50 \cdot 10^{-8}$
m^2 (GeV^2)	0.249	0.295	0.305	0.320	0.365	0.396	0.440	0.460	0.515
λ^2 (GeV^2)	0.0005	0.0002	0.00025	0.00015	0.0001	0.00009	0.00007	0.00005	0.00003
A (pn)	$1.65 \cdot 10^{-4}$	$3.09 \cdot 10^{-5}$	$1.02 \cdot 10^{-6}$	$1.14 \cdot 10^{-5}$	$2.46 \cdot 10^{-6}$	$9.44 \cdot 10^{-7}$	$2.2 \cdot 10^{-7}$	$1.16 \cdot 10^{-7}$	$9.4 \cdot 10^{-9}$

Asymmetry at $E_{lab} = 0.8$ GeV; distribution functions from Ref. 12.

TABLE 4.

m^2 (GeV^2)	0.280	0.285	0.292	0.300	0.320	0.4095	0.500	0.500
λ^2 (GeV^2)	0.0002	0.00025	0.0001	0.00015	0.00007	0.00009	0.00003	0.00005
$A(pp)$	$1.15 \cdot 10^{-5}$	$8.32 \cdot 10^{-6}$	$6.47 \cdot 10^{-6}$	$6.01 \cdot 10^{-6}$	$2.92 \cdot 10^{-7}$	$2.15 \cdot 10^{-7}$	$9 \cdot 10^{-9}$	$1.5 \cdot 10^{-8}$
m^2 (GeV^2)	0.250	0.250	0.264	0.275	0.309	0.320	0.330	0.365
λ^2 (GeV^2)	0.0002	0.00025	0,0001	0.00015	0.00009	0.00005	0.00007	0.00003
$A(pn)$	$2.46 \cdot 10^{-5}$	$3.02 \cdot 10^{-5}$	$1.48 \cdot 10^{-5}$	$1.36 \cdot 10^{-5}$	$3.67 \cdot 10^{-6}$	$2.3 \cdot 10^{-6}$	$2.40 \cdot 10^{-6}$	$5.42 \cdot 10^{-8}$

Asymmetry at $E_{lab} = 0.8$ GeV; distribution functions from Ref. 13.

TABLE 5.

m^2 (GeV^2)	0.65							
λ^2 (GeV^2)	0.5	0.05	0.001	0.0005	0.0001	0.00005	0.00001	0.000005
$\sigma_+ + \sigma_-$ (GeV^{-2})	$1.25 \cdot 10^{-2}$	0.34	28.96	58.36	292.98	586.05	2929.92	0
$\sigma_+ - \sigma_-$ (GeV^{-2})	$2.65 \cdot 10^{-4}$	$2.27 \cdot 10^{-4}$	$2.22 \cdot 10^{-4}$	$2.22 \cdot 10^{-4}$	$2.21 \cdot 10^{-4}$	$2.21 \cdot 10^{-4}$	$2.20 \cdot 10^{-4}$	0
m^2 (GeV^2)	0.75							
λ^2	0.5	0.05	0.001	0.0005	0.0001	0.00005	0.00001	0.000005
$\sigma_+ + \sigma_-$ GeV^{-2})	$2.1 \cdot 10^{-2}$	0.62	47.88	96.14	481.91	963.98	4820.19	9640.34
$\sigma_+ + \sigma_-$ (GeV^{-2})	$1.04 \cdot 10^{-3}$	$8.76 \cdot 10^{-4}$	$8.58 \cdot 10^{-4}$	$8.58 \cdot 10^{-4}$	$8.56 \cdot 10^{-4}$	$8.56 \cdot 10^{-4}$	$8.54 \cdot 10^{-4}$	$8.54 \cdot 10^{-4}$

Cross sections as functions of m^2 and λ^2 at $p_{lab} = 6$ GeV/c; distribution functions from Ref. 12.

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NARUŠENJE PARITETA U VISOKOENERTETSKOM PROTON-JEZGRA RASPRŠENJU I KVARK-PARTON MODEL

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Izračunato je narušenje pariteta u visokoenergetskom proton-jezgra raspršenju u partonskom modelu u kojem je struktura protona opisana protonskim distribucionim funkcijama s Q^2 zavisnošću. Eksperimentalni rezultati za raspršenje polariziranih protona na vodenoj meti mogu se reproducirati odgovarajućim parametrima modela.