Polarization effects in high-energy baryon reactions in QCD

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The polarization effects in $2\to 2$ processes involving baryons do not vanish with increasing energy. A version of QCD dual sum rules proposed here makes it possible to calculate the vertex of a nucleon with a pomeron.

It is widely believed that polarization effects in the scattering of hadrons consisting of massless u and d quarks "die out" with increasing energy. In the present letter we calculate the seed amplitudes for πN scattering as an example to demonstrate that this expectation is not justified in QCD.

We recall that the impact factor corresponding to a diagram of the type in Fig. 1a actually reduces to the corresponding form factor in Fig. 1b at high energies.[1,2] In order to clarify this question, it is therefore natural to begin with a study of the form factor in the corresponding kinematics. It is convenient to use the Sudakov variables $\tilde{p}_1 = p_1 - p_1^2 s^{-1} p_2$, $\tilde{p}_2 = p_2 - p_2^2 s^{-1} p_1$ for these calculations, writing all the vectors of the problem in the form $a = \alpha\tilde{p}_1 + \beta\tilde{p}_2 + a_\perp$, and writing $g^{\mu\nu}$ as $g^{\mu\nu} = g_\perp^{\mu\nu} + 2(\tilde{p}_1^\mu \tilde{p}_2^\nu + \tilde{p}_1^\nu \tilde{p}_2^\mu)/s$. In the leading approximation as $s\to\infty$ we can use the substitution $g^{\mu\nu}\to 2\tilde{p}_1^\nu\tilde{p}_2^\mu/s$ (here and below, the notation used for the momenta and the positions of the indices in the text correspond to Fig. 1). We can now reformulate the question in the following way: Is a helicity flip of the nucleon possible for the form factor $\langle N(p_2)|J^\mu|N(p_4)\rangle\tilde{p}_1^\mu$ in the massless limit for u and d quarks? At first glance, the answer would appear to be negative: Helicity is conserved. Nevertheless, helicity flip is possible, since quarks with fixed helicities which are not localized at one point go over into nucleons with both spin projections in the course of the hadronization.

We turn now to some quantitative estimates. To calculate the form factor introduced above, we use the correlation function with chiral currents, $I_1^{\lambda\lambda'}$, for which the nonperturbative corrections are small (15–20%) in comparison with the perturbative contribution over the entire range of $q_\perp^2$ (Ref. 3):

$$I_1^{\lambda\lambda'}(q_\perp^2) = i^2\int dx dy \exp(ip_2x + iqy) < 0|T\{\eta_L^\lambda(x)J_\mu(y)\bar{\eta}_L^{\lambda'}(0)\}|0 > \tilde{p}_1^\mu, \quad (1)$$

$$\eta_L^\lambda = (u_L^i C d_L^j)u_L^{k,\lambda}\epsilon^{i,jk}; \quad < 0|\eta_L^\lambda(0)|p >= f_N^L N_L^\lambda(p); \quad |f_N^L| \cong 0.38\cdot 10^{-2}\,\mathrm{GeV}^3.$$

The contribution of the nucleon to the spectral density of the correlation function $I_1^{\lambda\lambda'}$ is

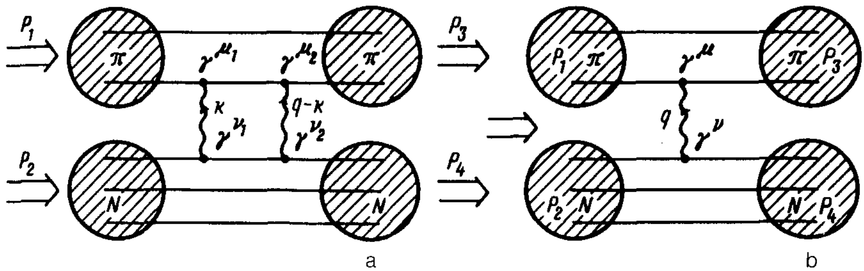

FIG. 1. $s=(p_1+p_2)^2$; $t=q_\perp^2\approx-\mathbf{q}_\perp^2$; $q=p_3-p_1$; $p_1+p_2=p_3+p_4$; $p_4\cong\hat{p}_2+\frac{p_2^2+\mathbf{q}_\perp^2}{s}\tilde{p}_1-q_\perp$.

$$|f_N|^2\pi^2\delta(s_1-M_N^2)\delta(s_2-M_N^2)[G_m(q_\perp^2)\{(1+\gamma_5)/2[\hat{p}_2\hat{\tilde{p}}_1\hat{p}_4+M_N^2\hat{\tilde{p}}_1]\}_{\lambda\lambda'}$$

$$-sF_2(q_\perp^2)M_N\{\hat{p}_2+\hat{p}_4\}_{\lambda\lambda'}]. \qquad (2)$$

To evaluate the electromagnetic form factors G_m and F_2, it is convenient to use the Lorentz structures $\varepsilon_{\mu\nu\alpha\beta}\tilde{p}_1^\mu\tilde{p}_2^\nu q_\perp^\alpha\gamma_{\lambda\lambda'}^\mu$, $\hat{q}_\perp$, and $\hat{p}_2$. Integrating the spectral density $\rho(s_1,s_2)$ with a weight $\exp(-(s_1+s_2)/m^2)$ from 0 to S_0 (S_0 is the nucleon duality interval, and m^2 is the Borel parameter), and ignoring the nonperturbative corrections, we find sum rules for the proton and neutron form factors. Since these expressions are rather lengthy, we will write out here only the sum rules for the magnetic form factors of the proton and the neutron:

$$G_m^{P,N}(q_\perp^2)=\frac{1}{A}\exp(2M_N^2/m^2)\beta^{P,N}\int_0^1 dx\mu^3(1-x)^2\exp(-2\tau/m^2);$$

$$\tau=q_\perp^2(1-x)/4x;$$

$$\mu^3=\frac{m^2}{8}[1-e^{-z}(1+z+\frac{z^2}{2})]\theta(z);\qquad z=2(S_0-\tau)/m^2;\ \ A=2^8\pi^4|f_N|^2/9; \quad (3)$$

$\beta^P=2e_u-e_d/2=3/2$; $\beta^N=2e_d-e_u/2=-1$; and $\theta(z)$ is the unit step function. Sum rules (3) yield values for the form factors $G_m^{P,N}(q_\perp^2)$ and $F_2^{P,N}(q_\perp^2)$ which agree well with experimental data to $q_\perp^2=3$–$5\ \mathrm{GeV}^2$. The implication is that this correlation function conveys the properties of the total wave function of the nucleon even at a perturbative level. The most important point for our purposes here is that the right-hand (theoretical) part of the sum rules in (3), which corresponds to the contribution of the form factor F_2 to the spectral density, does not vanish in the chiral limit even at the perturbative level, in accordance with the assertion above. A fit of sum rules (3) with $S_0 = 1.3$ GeV^2 reproduces the mass of a nucleon quite well: $M_N^2 = 0.85$–0.95 GeV^2. The value of S_0 agrees within 10% with the value of the duality interval for the two-point correlation function of the chiral currents which we are using.

We turn now to the scattering $\pi N \to \pi N$. We write the amplitudes for the process in the form $M = \hat{\bar{N}}(p_2)[T_1 s + \hat{\bar{p}}_1 T_2] N(p_4)$. The helicity amplitudes A^{++} and A^{+-} at high energies are expressed in terms of the amplitudes T_1 and T_2: $A^{+-} = \sqrt{-t}\, s T_1$; $A^{++} = s(2T_1 M_N + T_2)$. To calculate these helicity amplitudes we consider the correlation function

$$I_2^{\lambda\lambda'} = \int dx dy dz \exp(p_1 x + p_2 y - p_3 z) < 0|T\{\eta_L^\lambda(x) J_{\mu 5}(y) \bar{\eta}_L^{\lambda'}(z) J_{\nu 5}(0)\}|0 > \tilde{p}_2^\mu \tilde{p}_2^\nu.$$

Here $J_{\mu 5}$ is an interpolating current for the contribution of the π meson. Now proceeding as above, we use the dispersion relation in terms of the variables $p_1^2, p_2^2, p_3^2, p_4^2$ (Fig. 1). The contribution of the amplitudes in which we are interested to the spectral density of the correlation function $I_2^{\lambda\lambda'}$ is analogous to (2), with the substitutions $G_m \to T_{2j} F_2 \to -T_1$ and with a multiplication by the corresponding factor associated with the π-meson contribution. Evaluating the contributions corresponding to the diagrams in Figs. 1a and 2, and subtracting the continuum contribution from the theoretical part of the sum rules, we find sum rules for the helicity amplitudes. The very fact that the contribution of diagrams like that in Fig. 1a corresponds to a calculation of isoscalar form factors and the fact that $F_2^N + F_2^P$ is not zero guarantee that the amplitude T_1 will be different from zero. In other words, an amplitude with helicity flip makes a contribution to the differential cross section which does not decay with increasing energy. A calculation shows that A^{+-} is an order of magnitude smaller than A^{++}, but it increases more slowly with $|t|$ at $|t| \gg \mu^2$ than does the amplitude without helicity flip. We wish to stress that diagrams like that in Fig. 2 play an important role here. They become governing factors at sufficiently large values of the momentum transfer, as in reactions without helicity flip (see also Ref. 4). This example shows that in calculating scattering amplitudes by the approach of Ref. 2 it is necessary to consider the isoscalar form factor F_2 along with F_1, which is usually taken

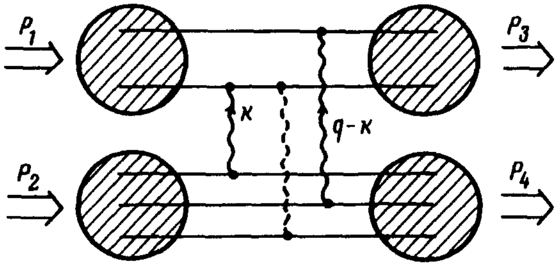

FIG. 2. Diagrams of this type cannot be reduced to a form factor; they are predominant at large values of $|t|$. The asymptotic behavior as $|t| \to \infty$ is determined by diagrams with an exchange of three gluons in the t channel. The third gluon, represented by the dashed line, appears in the next higher order of perturbation theory.

into consideration. In order to use the results in the region of small momentum transfer, we should also take account of the modification of the gluon propagator in the small-virtuality region.[5] We believe that the gluon component of the hadron wave function also plays an important role at small values of $|t|$. In the region $|t| \gtrsim 1\ \text{GeV}^2$, however, the sum rules must correctly convey the overall picture. Here we are not considering the higher orders of perturbation theory, but the results of Ref. 6 suggest that these higher orders could lead to only an increase in the cross sections with increasing energy. If so, our qualitative assertion regarding the role played by polarization effects would remain valid. In conclusion, we wish to stress that the mechanism discussed here for the flip of the baryon helicity is associated with the hadronization of massless quarks.

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