EXTRACTION OF SPIN OBSERVABLES IN BARYON-BARYON SCATTERING, SENSITIVE TO GLUON- AND QUARK-EXCHANGE EFFECTS

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(Received May 11, 2007)

Spin-dependent characteristics of baryon-baryon scattering are discussed, in which QCD dynamics is expected to display itself in most simply treatable regimes: 1) being dominated by semi-classical gluonic field without dynamical quarks; 2) as perturbative scattering. It is argued that peripheral semi-classical gluon exchange can make a major contribution to the pseudovector scattering amplitude, whereas one of the tensor amplitude components must be seeded by quark exchange, thus containing a hard scale and allowing for perturbative treatment. Expressions for absolute values of the promising amplitudes are derived in terms of double spin asymmetries in NN-scattering. Options for realization of corresponding polarization measurements in strange hyperon collisions with nucleons are analysed, along with consistency tests supplied by additional spin amplitudes emerging in non-identical particle collisions.

PACS: 11.80.Cr, 13.75.Cs, 13.85.Dz, 13.88.+e

1. INTRODUCTION

Accepting QCD as a fundamental theory of strong interactions, the ultimate goal is to obtain in terms of quark and gluon fields a spatial picture of all known hadrons, and find their dynamical susceptibilities (polarizabilities and instabilities to finite perturbations), that could serve as a ground for consistent calculations of the variety of hadronic processes. Systems, which may be simpler to try building such a detailed picture for, are mesons, especially those which contain heavy quarks (quarkonia). However, because quarkonium beams are not experimentally available (their ranges are typically of order millimeters), those particles can only be studied by their decays immediately after production. At that, in radiative and leptonic decays, supplying the cleanest and most direct information on quark distributions, that can be obtained only on the meson own mass scale, around a few GeV. Meanwhile, most detailed and diverse data relate to baryons - lepton scattering experiments provide knowledge of single quark distributions, and the subsequent solution of evolution equations offers incoherent gluon distributions at Bjørken x not too close to 1 and 0. But in order to probe spatial and temporal correlations between quarks and gluons in a nucleon, it is mandatory next to engage data on hadron-hadron scattering, at different momentum transfer and energy scales.

Qualitatively, the origin of empirical properties of baryon-baryon interactions allows for interpretation in terms of quark and gluon exchanges. Suggestions were made that NN-attraction can be described in terms of quark exchanges, similarly to hydrogen bonding in molecules due to electron exchange [1]. The repulsive core of NN-interaction of a few hundreds MeV in magnitude can be explained by effects of Fermi-degeneracy between constituent quarks. But in numerical accuracy, the quark model description of NN-scattering lags far behind the best-fit meson exchange models. The cause for greater success of meson models may be attributed to a weaker interaction between pions under non-resonance conditions, than that between quarks and gluons. That pion interaction with nucleons is strong, makes no further problem if one constructs a NN-potential[1] as a superposition of one-meson-exchange potentials [2], with the account, where necessary, of resonances in pion pairs, etc. [3] (see also [4]).

On the other hand, with the knowledge on the microscopic level that momentum transfer to an hadron can well result from exchange of gluons and quarks, despite that those particles do not exist in a free state, and besides that the quark distribution in baryons has a formidable spatial dimension, it becomes doubtful that local meson field theories are in principle capable of explaining *all* the properties of baryon-baryon interaction. In fact, at bringing the meson theory into correspondence with data, there exists a number of controversies, which stand unresolved over years.

1. In attempts to describe the spin-orbit part of interaction between nucleons at low energies via

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[1]The "potential", generically, is non-local, because it depends not only on positions, but also on momenta, however, that dependence is relatively small, and speaking of a "potential" is common.

exchange of known mesons, there has been no success to explain the large magnitude of the analyzing power A_y at $q \leq 70\,MeV/c$ (A_y-puzzle; see recent review [5]). The problem is that both π and 2π exchanges (including exchanges of ρ and ε viewed as 2π resonances) do not contribute so strongly to spin-orbit interaction, whereas the interaction through ω-exchange is too short-range and its coupling constant is bound by data on ω production. The physical origin of the given discrepancy may be due to substantial influence in the peripheral region of gluons, which are vector quanta. On the other hand, at highest energies (RHIC) inelastic *pp* cross-section makes up $10\,fm^2$, i. e. inelastic processes occur at impact parameters as large as $2fm$ between proton centers. At that, relying on the observed stable growth of cross-sections with energy, one may expect their further increase at least by a factor of $1.5-2$. That would already mean that inelastic processes occur intensely at interaction of nucleon peripheral regions. Because meson exchange can not provide the non-vanishing value of the cross-section at asymptotically high energies, since its probability falls off as a power law (even at vacuum quantum numbers - just due to existence of internal momentum distribution of constituents), such observations can give direct evidence of an essential role played by the gluon component in the region $1\,fm < r < 2\,fm$.

2. The meson core of inter-nucleon potential is not consistently defined. There is a discrepancy between the core width in the potential and that in the measured πNN formfactors - it is embarrassing to see that in meson potentials the core is about twice narrower than quark distribution in the nucleon as measured by direct EM and weak probes. It should also be recalled that the working assumption of a point-like nucleon is perhaps not reliable even in the limit $N_c \to \infty$, when it is correct for mesons [6]. In general, without making reference to the notion of constituent quarks, it is hard to explain the very πNN formfactors, steeply decreasing with the growth of the momentum transfer, the empirical law of nucleon-antinucleon pair suppression, etc. A detailed critique of meson exchange models may be found in [7].

Thus, without the account for spatial QCD structure of nucleons, description of their interactions meets felt difficulties both at large and at short distances, and moreover, in the relativistic domain. A logical approach for improvement of the correspondence with data might be to introduce QCD corrections into already approved meson exchange models. However, it is by now far from clear, how to build hybrid meson-chromodynamic models and in so doing to avoid double counting. Although some such attempts were made, they did not manage to refine general correspondence with data [8].

Perhaps, a more consistent, and yet not despairingly hardly realizable approach would be to apply pure chromodynamics, but only in those kinematic domains and for those variables, where meson exchange models do not offer an appropriate description. At that, one should look for maximal simplification of QCD equations under favorable kinematic conditions. Such conditions are the soft domain, where it is legitimate to use semi-classical Yang-Mills equations without dynamical quarks, or the hard domain, in which perturbation theory is applicable. It is natural to expect, that gluon exchange in the peripheral domain has the best chance to be semi-classical, whereas exchange by constituent quarks between baryons should introduce a relatively hard scale into the process, and so gives a chance for perturbative description.

The first objective, thus, is to extract a scattering amplitude, which is dominated by interaction in the peripheral domain and does not contain contributions from quark exchange. As we had noted in general, the peripheral pion exchange contributes to all possible forms of spin interaction but to spin-orbit. That means that one should consider a pseudovector scattering amplitude, which is most closely related to spin-orbit interaction. Procedure of extraction of the absolute value of this amplitude from scattering data will be described in Sec. 2.

Secondly, it is of importance to find amplitudes of processes, which necessarily require quark exchange, and thereby are not obscured by *soft* gluon exchange effects. Such amplitudes then might allow for perturbative treatment, though, possibly, with essential vertex renormalization, account for a spatially-inhomogeneous, dynamical condensate, etc. Note, that in general it is desirable for the process to involve a minimal number of degrees of freedom. In the case of hadron scattering, that means involving less quarks into violent dynamics, in order for the color confinement mechanism we poorly understand not to fully come to action. These conditions seem to be best satisfied by the nucleon charge exchange process, which will be discussed in Sec.3.

In Sec.4 we touch upon advantages gained for the named purposes with measurement of nucleon collisions with strange hyperons and discuss spin amplitudes which additionally arise in that case.

For simplicity of the presentation, in capacity of transition matrices between on-mass-shell spin states we will use Pauli matrices, though applications will be assumed in the relativistic domain. For the latter case the matrices may be understood as in [9], or it just could be noted, that upon transition to the 4-component covariant formalism [10], all relations remain qualitatively unchanged.

2. EXTRACTION OF GLUON EXCHANGE EFFECTS IN PERIPHERAL NN SCATTERING

Consider the matrix amplitude of two dynamically identical (within the unitary symmetry group) baryons of spin 1/2,

$$1 + 2 \to 3 + 4.$$

In the simplest experimental setting, such baryons are nucleons. At description of the process in the c.m.s. ($\mathbf{p}_1 + \mathbf{p}_2 = \mathbf{p}_3 + \mathbf{p}_4 = 0$), with the use of the orthogonal basis

$$\mathbf{q} = \mathbf{p}_3 - \mathbf{p}_1, \mathbf{r} = \mathbf{p}_3 + \mathbf{p}_1 = -\mathbf{p}_4 - \mathbf{p}_2, \mathbf{N} = [\mathbf{qr}],$$

parametrization of the matrix amplitude reads [11]

$$M = S_{31,42} I_{31} I_{42} + A^N_{31,42} \frac{1}{2} \left(\sigma^N_{31} I_{42} + I_{31} \sigma^N_{42} \right) + B^{NN}_{31,42} \sigma^N_{31} \sigma^N_{42} + B^{qq}_{31,42} \sigma^q_{31} \sigma^q_{42} + B^{rr}_{31,42} \sigma^r_{31} \sigma^r_{42}. \quad (1)$$

Here σ^α_{31} are Pauli matrices, operating from the spin space of particle 1 to the spin space of particle 3, and σ^β_{42} is the same for $2 \to 4$. The isospin indices are suppressed. The amplitude normalization will be supposed such that the differential cross-section averaged over initial and summed over final polarizations of particles expresses as

$$\left\langle \frac{d\sigma}{dt} \right\rangle = \frac{1}{2} Sp \frac{1}{2} Sp M M^\dagger = \left| S_{31,42} \right|^2 + \frac{1}{2} \left| A^N_{31,42} \right|^2 + \left| B^{NN}_{31,42} \right|^2 + \left| B^{qq}_{31,42} \right|^2 + \left| B^{rr}_{31,42} \right|^2 . \quad (2)$$

The representation (1) may be called covariant, in contrast to, e. g., treatment of the amplitude in a fixed spin basis. It is the covariant representation for the amplitude, which proves convenient for extraction of exchanges with specific quantum numbers.

In order to separate in the given scattering process effects, which might reasonably be regarded as dominated by gluon exchange, let us approach the nucleon from large distances and consider the gradual actuation of contributions to amplitude (1) from known mesons. The longest range is ascribed to one-pion exchange, which contributes mainly to $B^{qq}_{31,42}$ and, upon iteration, to $S_{31,42}$ [12]. At a closer approach to the nucleon, correlated 2π exchanges become noticeable, which in the simplest treatment [13], yield

$$A^N_{31,42} \approx 0, \quad B^{qq}_{31,42} \approx 0, \quad S_{31,42} \approx B^{NN}_{31,42} \approx B^{rr}_{31,42} \neq 0.$$

In other versions of 2π-exchange calculations (see their comparison in [14]) departures from that scenario are possible, but at any rate, as translated to potentials, the pion-theoretic spin-orbit interaction may be regarded as vanishing at $r > 1\, fm$ (whereas the scalar and tensor potentials retain considerable magnitude down to as far as $r \approx 1.5 - 2\, fm$). Thus, if the gluon cloud does extend to distances $\sim 2\, fm$, its effect will be most unambiguously discernible on the pion cloud background in the pseudovector amplitude[2] $A^N_{31,42}$.

Separating the scattering at impact parameters $1\, fm < r < 2\, fm$ from the region $r < 1 fm$ would be sufficiently feasible in case if the scattering proceeded as a semi-classical deflection in a strong potential field. However, NN-interaction in the region $1 - 2\, fm$, where the nuclear potential has values of the order $20\, MeV$, may be regarded as strong only at energies lower than that scale, whereas for spatial resolution of $1\, fm$ distances one needs momenta $> 150\, MeV/c$, i. e. nucleon energies $> 10\, MeV$. So, it appears like there is no way to avoid quantum effects and intermixture of central and peripheral region contributions. Nonetheless, when considering diffractive scattering at $E > 1\, GeV$, the contribution to the spatial profile of elastic scattering from the impact parameter region $b < 1\, fm$ may prove to be rather small due to high inelasticity (for elastic amplitude equivalent to opaqueness). Also, specifically in the case of the amplitude $A^N_{31,42}$, with the account for quark structure of the nucleon, it may happen that the spin-orbit interaction in the region between quarks in average is small. That conjecture finds confirmation, for instance, in the Isgur-Karl constituent quark model [15], [16], in which the observed spectrum of P-baryons is nicely reproduced with the neglect of LS-interaction between quarks.

The observable characteristic of $A^N_{31,42}$ in the case of $NN \to NN$ is usually thought to be single-spin asymmetry [3]

$$P_N = \frac{d\sigma_\uparrow - d\sigma_\downarrow}{d\sigma_\uparrow + d\sigma_\downarrow} = \frac{\frac{1}{2} Sp \frac{1}{2} Sp M \sigma^N_{11} M^\dagger}{\frac{1}{2} Sp \frac{1}{2} Sp M M^\dagger}, \quad (3)$$

where the spin projection direction is supposed to be $\mathbf{N}$. However, through the spin amplitudes present in (1) the asymmetry P_N expresses as[4]

$$P_N = \left\langle \frac{d\sigma}{dt} \right\rangle^{-1} Re \left(S_{31,42} + B^{NN}_{31,42} \right) A^{N*}_{31,42},$$

and despite being proportional to $A^N_{31,42}$, it also contains other amplitudes. Hence, this asymmetry quantitatively tells little about $A^N_{31,42}$ [17].

A non-trivial task, therefore, is to measure $A^N_{31,42}$ separately from other spin amplitudes. In fact, one may expect only to determine $\left| A^N_{31,42} \right|^2$ and the complex phase of $A^N_{31,42}$ relative to other amplitudes. The next grade in measurement complexity is double-spin asymmetries. Those are

$$C_{\beta\alpha} = \left\langle \frac{d\sigma}{dt} \right\rangle^{-1} \frac{1}{2} Sp \frac{1}{2} Sp M \sigma^\beta_{22} \sigma^\alpha_{11} M^\dagger,$$

$$D_{\beta\alpha} = \left\langle \frac{d\sigma}{dt} \right\rangle^{-1} \frac{1}{2} Sp \frac{1}{2} Sp \sigma^\beta_{33} M \sigma^\alpha_{11} M^\dagger,$$

[2] Occasionally, we will refer to $A^N_{31,42}$ as to pseudovector amplitude, though it is understood in fact to be just the only component of the pseudovector amplitude, which is not bound to be zero due to some symmetry reasons or other.

[3] The quantity P_N also (by Madison convention [18]) is denoted as A_y; thereunder it was mentioned in the Introduction.

[4] For the Hermitean conjugate matrix $M^\dagger$ the notation concordant with (1) is
$M = S^*_{31,42} I_{13} I_{24} + A^{N*}_{31,42} \frac{1}{2} \left(\sigma^N_{13} I_{24} + I_{13} \sigma^N_{24} \right) + B^{NN*}_{31,42} \sigma^N_{13} \sigma^N_{24} + B^{qq*}_{31,42} \sigma^q_{13} \sigma^q_{24} + B^{rr*}_{31,42} \sigma^r_{13} \sigma^r_{24}$,
with particle order in matrix subscripts inverted.

$$K_{\beta\alpha} = \left\langle \frac{d\sigma}{dt} \right\rangle^{-1} \frac{1}{2} Sp \frac{1}{2} Sp \sigma_{33}^{\beta} M \sigma_{22}^{\alpha} M^{\dagger},$$

spin correlation parameters (otherwise denoted as[5] $A_{\beta\alpha}$); depolarization and spin rotation parameters, and polarization transfer parameters. In the basis (1), non-zero are 9 diagonal asymmetries

$$\begin{aligned}
C_{NN} \left\langle \frac{d\sigma}{dt} \right\rangle &= 2Re\left(S_{31,42} B_{31,42}^{NN*}\right) + \frac{1}{2}\left|A_{31,42}^{N}\right|^2 - 2Re\left(B_{31,42}^{qq} B_{31,42}^{rr*}\right), \\
C_{qq} \left\langle \frac{d\sigma}{dt} \right\rangle &= 2Re\left(S_{31,42} B_{31,42}^{qq*}\right) - 2Re\left(B_{31,42}^{NN} B_{31,42}^{rr*}\right), \\
C_{rr} \left\langle \frac{d\sigma}{dt} \right\rangle &= 2Re\left(S_{31,42} B_{31,42}^{rr*}\right) - 2Re\left(B_{31,42}^{NN} B_{31,42}^{qq*}\right), \\
D_{NN} \left\langle \frac{d\sigma}{dt} \right\rangle &= \left|S_{31,42}\right|^2 + \frac{1}{2}\left|A_{31,42}^{N}\right|^2 + \left|B_{31,42}^{NN}\right|^2 - \left|B_{31,42}^{qq}\right|^2 - \left|B_{31,42}^{rr}\right|^2, \\
D_{qq} \left\langle \frac{d\sigma}{dt} \right\rangle &= \left|S_{31,42}\right|^2 - \left|B_{31,42}^{NN}\right|^2 + \left|B_{31,42}^{qq}\right|^2 - \left|B_{31,42}^{rr}\right|^2, \qquad (4) \\
D_{rr} \left\langle \frac{d\sigma}{dt} \right\rangle &= \left|S_{31,42}\right|^2 - \left|B_{31,42}^{NN}\right|^2 - \left|B_{31,42}^{qq}\right|^2 + \left|B_{31,42}^{rr}\right|^2, \\
K_{NN} \left\langle \frac{d\sigma}{dt} \right\rangle &= 2Re\left(S_{31,42} B_{31,42}^{NN*}\right) + \frac{1}{2}\left|A_{31,42}^{N}\right|^2 + 2Re\left(B_{31,42}^{qq} B_{31,42}^{rr*}\right), \\
K_{qq} \left\langle \frac{d\sigma}{dt} \right\rangle &= 2Re\left(S_{31,42} B_{31,42}^{qq*}\right) + 2Re\left(B_{31,42}^{NN} B_{31,42}^{rr*}\right), \\
K_{rr} \left\langle \frac{d\sigma}{dt} \right\rangle &= 2Re\left(S_{31,42} B_{31,42}^{rr*}\right) + 2Re\left(B_{31,42}^{NN} B_{31,42}^{qq*}\right),
\end{aligned}$$

and 3 antisymmetric non-diagonal asymmetries

$$\begin{aligned}
D_{qr}\langle\frac{d\sigma}{dt}\rangle &= -D_{rq}\langle\frac{d\sigma}{dt}\rangle = Im\left[A_{31,42}^{N}(S_{31,42}^{*} - B_{31,42}^{NN*})\right]\varepsilon^{qrN}, \\
C_{qr}\langle\frac{d\sigma}{dt}\rangle &= -C_{rq}\langle\frac{d\sigma}{dt}\rangle = Im\left[A_{31,42}^{N}(B_{31,42}^{qq*} - B_{31,42}^{rr*})\right]\varepsilon^{qrN}, \qquad (5) \\
K_{qr}\langle\frac{d\sigma}{dt}\rangle &= -K_{rq}\langle\frac{d\sigma}{dt}\rangle = Im\left[A_{31,42}^{N}(B_{31,42}^{qq*} + B_{31,42}^{rr*})\right]\varepsilon^{qrN}.
\end{aligned}$$

The 9 relations (4) may be regarded as a complete system of equations for determination of 5 complex numbers $S_{31,42}$, $A_{31,42}^{N}$, $B_{31,42}^{NN}$, $B_{31,42}^{qq}$, $B_{31,42}^{rr}$ up to a common for them phase factor. This nonlinear system of equations can be explicitly solved. Firstly, through the measured asymmetries it is possible to express 3 quantities

$$4\left|B_{31,42}^{qq}\right|^2 = \left\langle \frac{d\sigma}{dt} \right\rangle (1 - D_{NN} + D_{qq} - D_{rr}), \qquad (6)$$

$$4\left|B_{31,42}^{rr}\right|^2 = \left\langle \frac{d\sigma}{dt} \right\rangle (1 - D_{NN} - D_{qq} + D_{rr}), \qquad (7)$$

$$4Re\left(B_{31,42}^{qq} B_{31,42}^{rr*}\right) = \left\langle \frac{d\sigma}{dt} \right\rangle (K_{NN} - C_{NN}), \qquad (8)$$

related to the pair of complex amplitudes $B_{31,42}^{qq}$, $B_{31,42}^{rr}$. Those amplitudes may be used as a basis set in the 2d Euclidean space, to which the complex plane turns upon adoption in it a scalar product $(a, b) = Re\,(a^{*}b) = ReaReb + ImaImb$. Then, quantities (6-8) serve as Gram matrix elements in the basis formed by $B_{31,42}^{qq}$ and $B_{31,42}^{rr}$. Next, projections of $S_{31,42}$ on $B_{31,42}^{qq}$ and $B_{31,42}^{rr}$ can be expressed as

$$4Re\left(S_{31,42} B_{31,42}^{qq*}\right) = \left\langle \frac{d\sigma}{dt} \right\rangle (K_{qq} + C_{qq}),$$

$$4Re\left(S_{31,42} B_{31,42}^{rr*}\right) = \left\langle \frac{d\sigma}{dt} \right\rangle (K_{rr} + C_{rr}),$$

and thereupon $\left|S_{31,42}\right|^2$ is reconstructed by formula

$$\begin{aligned}
\left|S_{31,42}\right|^2 \quad = \quad & \Big\{\left[Re\left(S_{31,42} B_{31,42}^{qq*}\right)\right]^2 \left|B_{31,42}^{rr}\right|^2 + \left[Re\left(S_{31,42} B_{31,42}^{rr*}\right)\right]^2 \left|B_{31,42}^{qq}\right|^2 \\
& -2Re\left(S_{31,42} B_{31,42}^{rr*}\right) Re\left(S_{31,42} B_{31,42}^{qq*}\right) Re\left(B_{31,42}^{qq} B_{31,42}^{rr*}\right)\Big\} \qquad (9) \\
& \cdot \Big\{\left|B_{31,42}^{rr}\right|^2 \left|B_{31,42}^{qq}\right|^2 - \left[Re\left(B_{31,42}^{qq} B_{31,42}^{rr*}\right)\right]^2\Big\}^{-1}.
\end{aligned}$$

[5]To be precise, the initial state parameters are denoted as $A_{\beta\alpha}$, and final state ones - as $C_{\beta\alpha}$. However, under T-invariance conditions those quantities coincide, so we shall use a more convenient notation $C_{\beta\alpha}$.

Finally, $\left|A^N_{31,42}\right|^2$ can be expressed from the equation

$$4\left|S_{31,42}\right|^2 + \left|A^N_{31,42}\right|^2 = \left\langle \frac{d\sigma}{dt} \right\rangle (1 + D_{NN} + D_{qq} + D_{rr})$$

by substitution of (9). As a result, we arrive to the expression

$$\begin{aligned}
\left|A^N_{31,42}\right|^2 \left\langle \frac{d\sigma}{dt} \right\rangle^{-1} &= 1 + D_{NN} + D_{qq} + D_{rr} \\
&- \Big\{ (K_{qq} + C_{qq})^2 (1 - D_{NN} - D_{qq} + D_{rr}) + (K_{rr} + C_{rr})^2 (1 - D_{NN} + D_{qq} - D_{rr}) \\
&\quad -2 (K_{rr} + C_{rr}) (K_{qq} + C_{qq}) (K_{NN} - C_{NN}) \} \\
&\times \Big\{ (1 - D_{NN})^2 - (D_{qq} - D_{rr})^2 - (K_{NN} - C_{NN})^2 \Big\}^{-1} . \qquad (10)
\end{aligned}$$

This expression employs all 9 diagonal double-spin asymmetries and the non-polarized differential cross-section. Note that at small q the $\mathbf{r}$ direction is equivalent to commonly used in experiments direction l of initial particle collision, or final particle emergence, and direction $\mathbf{q}$ is equivalent to the sideways direction s in the scattering plane, with respect to momentum of either of the particles.

Having reconstructed the amplitude $A^N_{31,42}(q)$ absolute value in the diffractive domain $q < 700\,MeV/c$, one can subsequently try to estimate its absolute phase as well, exploiting analyticity properties and the behavior of single-spin asymmetry (3). If successful, it would be worth further to make transformation from momentum transfers to impact parameters, in order to check the conjecture that there is a suppression in the central region. Should it be the case, and the profile $A^N_{31,42}(b)$ of amplitude $A^N_{31,42}(q)$ is built up in the region $1\,fm < b < 2\,fm$, then calculation of amplitude $A^N_{31,42}$ behavior on the basis of gluon exchange without that of quarks may be sensible. Once more it should be emphasized that for applicability of the impact parameter representation the collision must proceed with velocity greater then those of mechanical oscillations and color circulation in the nucleon. To this end, perhaps, it suffices to have $E > 1\,GeV$.

3. AMPLITUDES DETERMINED BY QUARK EXCHANGE

Turning to the problem of separation of amplitudes governed by quark exchange, the first point that needs to be clarified in this business is which quarks had actually participated in exchange. Baryon beams available in experiments may contain only quarks of flavors u, d and s, so in colliding baryons some identical quarks are always present. For a flavor exchange reaction, which by itself requires exchange of different quarks, it remains to secure absence of additional exchange of identical quarks. For elastic scattering, amplitudes are needed which may differ from zero only if exchange of identical quarks had occurred (the identity of exchanged quarks here is required by the elasticity).

In the first case the most practical variant is to consider nucleon charge exchange $np \to pn$ reaction in the forward direction at high energy. The non-polarized differential cross-section for this process in the vicinity of 0° direction (that is 180° for elastic scattering $np \to np$) features a peak fittable by a sum of two Gaussians [19, 20],

$$\left\langle \frac{d\sigma}{dt} \right\rangle \approx \frac{const}{s^2} \left(e^{-50|t|/GeV^2} + 0.8 e^{-4.5|t|/GeV^2} \right).$$

Note here energy dependence $d\sigma/dt \propto 1/s^2$, whereas at other fixed angles the differential cross-section falls off as $d\sigma/dt \propto 1/s^{10}$.

Historically, the observation that the width of the first peak is $GeV^2/50 \sim m_\pi^2$ inspired modeling of this process in terms of one-pion exchange. At practice, however, OPE gives in the backward direction zero instead of a peak, so a pion plus pomeron exchange has rather to be considered, to produce a cut with special interference conditions [21]. (And still, as [22] notes, approach of [21] fails to reproduce behavior of all *polarized* observables).

From the standpoint of QCD, one may imagine the mechanism of the peak formation, in which a quark from the first nucleon upon knocking a quark in the second nucleon substitutes it in the very same state (except flavor), whereas the knocked quark, vice versa, is sent into the freed position of the quark in the first nucleon. Since no color neutralization is required thereat, no $q\bar{q}$ pair needs to be created, so, most of the probability can be retained within the 2-hadron channel. The differential cross-section of two relativistic quark scattering to 180° through one-gluon exchange is proportional to $1/s^2$, thus no immediate contradiction with the experiment emerges. Manifestation of some other, harder scale may be attributed to collective effects between constituent quarks. As a whole, the given process should be sensitive to quark wavefunction in a nucleon.

In the case of elastic pp scattering the only amplitude, which is obliged to be zero in the event of exchange exceptionally by gluons, is $B^{qq}_{31,42}$. Indeed, if a diagram of two-fermion scattering can be cut only through lines of vector quanta, T-invariance forbids appearance in the total amplitude associated with each half of the diagram of matrices σ^q. To get assured that the difference of $B^{qq}_{31,42}$ from zero is caused by exchange of only one quark pair, it is sufficient, again, to consider scattering at high energy to small angles.

It would be interesting to compare the amplitude $B^{qq}_{31,42}$ in elastic scattering $pp \to pp$, the amplitude $B^{qq}_{31,42}$ in scattering $np \to np$, and $B^{qq}_{31,42}$ in the reaction $np \to pn$. Naive quark counting shows that in $pp \to pp$ scattering there are 5 possibilities of exchange by a pair of identical quarks, whereas in $pn \to pn$ processes such possibilities are 4, and in $np \to pn$ there are 4 possibilities of exchange $d(n) \leftrightarrow u(p)$. One has thus two points to check: 1). whether forward $\left|B^{qq}_{31,42}(\theta)\right|^2$ for $pp \to pp$ and $pn \to pn$ are indeed in the ratio 5:4, and 2). does forward $B^{qq}_{31,42}$ for $pp \to pp$ and $pn \to pn$, if indeed governed by quark exchange, develop a small peak of width $\sim m_\pi^2$, similarly to what is known for the $np \to pn$ process?

Realization of the suggested program is hindered by the circumstance that, as formula (6)[6]

$$\left|B^{qq}_{31,42}\right|^2 = \frac{1}{4}\left\langle\frac{d\sigma}{dt}\right\rangle(1 - D_{NN} + D_{qq} - D_{rr}) \equiv \frac{1}{4}\left\langle\frac{d\sigma}{dt}\right\rangle(1 - D_{NN} + D_{ss} - D_{ll}) \qquad (11)$$

indicates, for determination of $\left|B^{qq}_{31,42}\right|^2$, final nucleon polarization measurements are required. By now, such data at small angles remain too scarce (in contrast to differential cross-sections, analyzing powers [23] and initial-state double-spin correlations). Customary complications of re-scattering measurements after scattering to small angles are aggravated by the fact that as $\theta \to 0$, each of the three parameters D_{NN} , D_{ss} , D_{ll} gets within 10% from 1, and it is this deviation which needs to be measured. The highest energy where all D parameters are available, at small scattering angles, is $0.8\,GeV$ $(s - 4m^2 = 1.5\,(GeV/c)^2)$. Combining data [24] for unpolarized $pp \to pp$ differential cross-sections with data [25], [26] for D parameters, we can reconstruct $\left|B^{qq}_{31,42}\right|^2$. Data [27], [26] allow to do the same for $pn \to pn$. Finally, for $np \to pn$ quantity $\left|B^{qq}_{31,42}\right|^2$ can also be obtained from data [28], [29], [30], if in 11 $D_{\beta\alpha}$ are substituted by the spin-transfer parameters $K_{\beta\alpha}$. The results together are shown in Fig. 1.

The obtained dependences make a two-fold impression. On the one hand, values for $pp \to pp$ and $pn \to pn$ are uniformly small, as it could be expected based on their origin from quark scattering to large angles. The ratio for processes $pp \to pp$ and $pn \to pn$ is close to 5/4, though under significant errors (especially in processes involving neutrons) and small number of points in the considered domain of small angles, it is impossible to claim a definite correspondence.

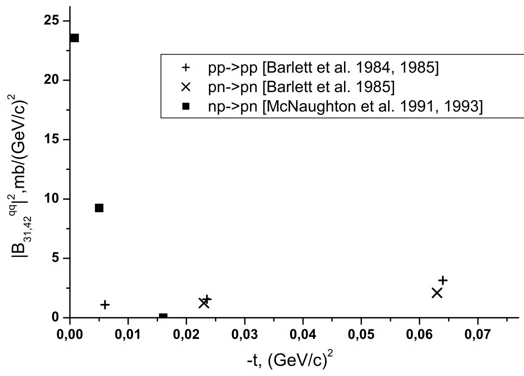

Fig.1. *Behavior of* $\left|B^{qq}_{31,42}\right|^2$ *for pp and pn elastic scattering, and for* $np \to pn$*, at* $0.8\,GeV$

On the other hand, in behavior of $\left|B^{qq}_{31,42}\right|^2$ no peak is seen on the scale $|t| \sim 0.01(GeV/c)^2$, although for $np \to pn$ such a peak is present, and even quite large, reaching one third from the unpolarized cross-section, though it in principle can not be larger then its half, since $B^{NN}_{31,41} \underset{\theta\to 0}{\approx} B^{qq}_{31,41}$ must make an equal contribution.

The cause for the strong difference in behavior of $\left|B^{qq}_{31,42}\right|^2$ for $pp \to pp$ and $np \to pn$ about $\theta \to 0$ is not clear. Meanwhile, though with a small likelihood, it can not be excluded that the leftmost point for $pp \to pp$ is inaccurate, e. g., due to a fallacy brought in by phase shift analysis predictions. In updated data [26], as compared to [25], that point is not present.

4. ADDITIONAL AMPLITUDES IN HYPERON-NUCLEON COLLISIONS

For measurement of baryon polarization in the final state a valuable advantage may come from observations of nucleon collisions with strange baryons $(YN \to YN)$, due to the fact that hyperon final polarizations are detectable by kinematics of their nonleptonic weak decays[7] $\Lambda \to N\pi$, $\Sigma^+ \to N\pi$. The lifetime for hyperons is $\sim 10^{-10}s$, so their range at $E \geq 1GeV$ exceeds $4cm$, which is sufficient for conduction of scattering experiments, though small angle measurements are difficult. Presently, differential cross-sections only to angles $\theta > 20^\circ$ were measured. At high energies, with the use of gas targets, there is potential for improving the situation [31]. Should measurements of hyperon diffractive scattering become practical, determination of quantity $\left|B^{qq}_{31,42}\right|^2$ and its comparison with the counter-

[6] Passage from the first line of (11), which is formula (6) to the second line, which implements directly measured quantities D_{ss}, D_{ll}, related to initial and final particle momenta directions, instead of a universal frame, proves to be exact, not just approximate.

[7] Transverse polarization of hyperons is easily measurable by the azimuthal asymmetry of pion emergence. The longitudinal component of polarization is harder to detect, but with the use of Dalitz diagrams, is also possible.

part in $pp \to pp$ will be of primary interest. Besides that, YN-scattering is interesting by additional amplitudes, which emerge in it due to hypercharge symmetry violation. Those amplitudes will be discussed hereafter.

In the process $YN \to YN$ such an additional amplitude comes as

$$A^{[N]}_{31,42}\frac{1}{2}\left(\sigma^N_{31}I_{42} - I_{31}\sigma^N_{42}\right). \tag{12}$$

This amplitude may be regarded as characteristic of spin-orbit dependence on quark mass. Analogously to what has been argued about $A^N_{31,42}$, amplitude $A^{[N]}_{31,42}$ as well may be subject to essential cancelations in the central region.

Representative of the order of magnitude of amplitude (12) as compared to $A^N_{31,42}$ is the ratio

$$\frac{2Re\left[\left(S_{31,42} - B^{NN}_{31,42}\right)A^{[N]*}_{31,42}\right]}{2Re\left[\left(S_{31,42} - B^{NN}_{31,42}\right)A^{N*}_{31,42}\right]} = \frac{\frac{1}{2}Sp\frac{1}{2}SpM\sigma^N_{11}M^\dagger - \frac{1}{2}Sp\frac{1}{2}SpM\sigma^N_{22}M^\dagger}{\frac{1}{2}Sp\frac{1}{2}SpM\sigma^N_{11}M^\dagger + \frac{1}{2}Sp\frac{1}{2}SpM\sigma^N_{22}M^\dagger}. \tag{13}$$

This ratio is convenient for its independence on the non-polarized cross-section magnitude, which is usually the source of significant experimental errors. Since we expect $A^{[N]}_{31,42}$ to be small, it does not essentially influence formulas (6), (10). Reproduction of the correct order of magnitude of the ratio (13) may serve as a test of correctness of the spin-orbit interaction calculation on the basis of quark models.

Exchange processes can also be studied in YN collisions[8], even without measurement of the hyperon polarization. It is natural to expect that in reaction $\Sigma^+p \to p\Sigma^+$, just like in $np \to pn$, a forward peak must develop. However, between those two cases two differences are to be minded. First of all, in the process $\Sigma^+p \to p\Sigma^+$ only s quark from Σ^+-hyperon and d-quark from proton may be engaged into exchange, so there are no Fermi-degeneracy effects in this case, and hence a harder scale in the peak shape may disappear. Secondly, if a baryon loses s-quark, and a d-quark comes in its place, the substitution is not exact, so there must occur some rearrangement of the wave function. Thus, examination of the difference between peaks in reactions $np \to pn$ and $\Sigma^+p \to p\Sigma^+$ can provide comparison between wave functions of nucleons and hyperons.

Further it may be noted that absolute normalization of cross-sections, necessary for the processes $np \to pn$ and $\Sigma^+p \to p\Sigma^+$ poses a formidable experimental problem, but that can be avoided, if for the same collisions, for example, Λp, different channels of charge exchange are simultaneously observed, such as $\Lambda p \to \Sigma^+n$ and $\Lambda p \to n\Sigma^+$. The first channel here must be similar to $np \to pn$, whereas the second - to $YN \to NY$. Therefore, in the forward direction peaks in Σ^+ and n particle distributions, originating form Λp collisions, should exhibit difference.

For processes of the mentioned type, when all 4 initial and final particles are different, it should be emphasized that the matrix amplitude represents sum of (1) and (12) only if in particle numeration one tracks strangeness, i. e. particles 1 and 3 both are hyperons. Dynamically, though, it may be convenient rather to track two quarks of the three. In the latter case, in the place of (12) there must stand an amplitude of the form[9]

$$B^{[qr]}_{31,42}\frac{1}{2}\left(\sigma^q_{31}\sigma^r_{42} - \sigma^r_{31}\sigma^q_{42}\right), \tag{14}$$

as is obvious from the spin crossing relation

$$iA^{[N]}_{31,42} = \varepsilon^{Nqr}B^{[qr]}_{41,42},$$

derived from Fierz identities for the $SU(2) \otimes SU(2)$ group. The corresponding difference of single-spin asymmetries reads as

$$2Im\left[\left(B^{qq}_{31,42} + B^{rr}_{31,42}\right)B^{[qr]*}_{41,42}\right]\varepsilon^{Nrq} = \frac{1}{2}Sp\frac{1}{2}SpM\sigma^N_{11}M^\dagger - \frac{1}{2}Sp\frac{1}{2}SpM\sigma^N_{22}M^\dagger,$$

$$\left(A^{[N]}_{31,42} = 0\right).$$

As a relative measure for amplitude $B^{[qr]}_{41,42}$ a ratio

$$\frac{2Im\left[\left(B^{qq}_{31,42} + B^{rr}_{31,42}\right)B^{[qr]*}_{41,42}\right]\varepsilon^{Nrq}}{2\left(\left|B^{qq}_{31,42}\right|^2 - \left|B^{rr}_{31,42}\right|^2\right)} = \frac{\frac{1}{2}Sp\frac{1}{2}SpM\sigma^N_{11}M^\dagger - \frac{1}{2}Sp\frac{1}{2}SpM\sigma^N_{22}M^\dagger}{\frac{1}{2}Sp\frac{1}{2}Sp\sigma^q_{33}M\sigma^q_{11}M^\dagger - \frac{1}{2}Sp\frac{1}{2}Sp\sigma^r_{33}M\sigma^r_{11}M^\dagger}. \tag{15}$$

can be considered. If we are poised to describe basic, diagonal tensor spin amplitudes based on pQCD and baryon wave functions, the framework should be able to reproduce the correct order of magnitude of the non-diagonal amplitude $B^{[qr]}_{41,42}$ as well.

5. SUMMARY

In this paper it was argued that when issuing from covariant representation (1) for the scattering amplitude, the amplitude $A^N_{31,42}(q)$ at relativistic energies and small scattering angles is dominated by gluon exchange contributions, and, presumably, from the peripheral region. The amplitude $B^{qq}_{31,42}(q)$, also at relativistic energies and small scattering angles, is seeded by quark exchange processes. As was discussed, those two amplitudes have best chances to be treatable by simplest approximations of QCD.

[8]Experiments of that kind are to be conducted at colliding beams, in order for hyperons after head-on collisions not to stop and decay too close to the scattering point.

[9]If the isospin symmetry violation is taken into account, too, as for example it should be in the region of Coulomb-nuclear interference, then both amplitudes ... and ... must be present, so in total there are 7 amplitudes, not 6.

For determination of absolute values of the mentioned complex amplitudes, expressions were derived, which contain the experimentally observable double-spin asymmetries of NN scattering process. At that, necessary ingredients are the depolarization and spin rotation parameters D_{NN}, D_{qq}, D_{rr}, which require measurement of polarization of one of the final nucleons. Currently, in connection with the development of nuclear polarized beam and target techniques, the main emphasis in experiments is laid on measurements double-spin correlations between initial particles. Those correlation show behavior, which still receives no satisfactory explanation in field-theoretic terms [32], and that may be the reason why more laborious measurements of final polarizations are not considered as urgent. Let us stress, however, that of particular value would be measurements which might be directly converted to indications concerning profile shapes or interaction amplitudes. Such unambiguous indications can be gained only upon the knowledge of spin correlations between initial and final particles.

Meanwhile, in the absence of sufficient set of polarization data, investigations of quark exchange processes can be carried out using non-polarized differential cross-sections of charge exchange processes, such as $np \to pn$. There are good reasons to hope that in processes of that kind the account for final state quark interactions may be made with the use of baryon wave functions, between which a perturbative QCD matrix element of quark exchange is evaluated. At practice, however, the question remains open, whether it is possible to use for this goal wave functions of potential or of bag models, which are essentially limiting cases corresponding to high and low gluonic transparency within a nucleon.

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ВЫДЕЛЕНИЕ СПИНОВЫХ НАБЛЮДАЕМЫХ В БАРИОН-БАРИОННОМ РАССЕЯНИИ, ЧУВСТВИТЕЛЬНЫХ К ЭФФЕКТАМ ОБМЕНА ГЛЮОНАМИ И ОБМЕНА КВАРКАМИ

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Обсуждаются спиновые характеристики барион-барионного рассеяния, в которых ожидается наименее замаскированное проявление КХД степеней свободы: 1) в квазиклассическом глюонном режиме без динамических кварков; 2) в пертурбативном режиме. Показано, что квазиклассический глюонный обмен наиболее отчетливо проявляется в псевдовекторной амплитуде рассеяния, тогда как тензорные спиновые амплитуды позволяют выделить процессы обмена кварками, содержащие относительно жесткий импульсный масштаб. Предложены выражения для абсолютных величин обсуждаемых амплитуд через двуспиновые асимметрии в упругом NN-рассеянии. Обсуждаются варианты проведения соответствующих поляризационных измерений в столкновениях странных гиперонов с нуклонами, и информация, предоставляемая появляющимися в этих процессах дополнительными амплитудами.

ВИДІЛЕННЯ СПІНОВИХ СПОСТЕРЕЖНИХ У БАРІОН-БАРІОННОМУ РОЗСІЯННІ, ЧУТЛИВИХ ДО ЕФЕКТІВ ОБМІНУ ГЛЮОНАМИ ТА КВАРКАМИ

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Обговорюються спінові характеристики баріон-баріонного розсіяння, в яких очікуються найменш замасковані прояви КХД ступенів свободи: 1) у квазикласичному глюонному режимі без динамічних кварків; 2) в пертурбативному режимі. Показано, що квазикласичний глюонний обмін найбільш виразно проявляється у псевдовекторній амплітуді розсіяння, тоді як тензорні спінові амплітуди дозволяють виділити процеси обміну кварками, що привносять відносно жорсткий імпульсний масштаб. Запропоновано вирази для абсолютних величин амплітуд, що обговорюються, через двоспінові асиметрії в пружному NN-розсіянні. Обговорюються можливості проведення відповідних поляризаційних вимірів у зіткненнях дивних гіперонів з нуклонами, а також інформація, що надається додатковими амплітудами у цих процесах.