Few-body aspects of non-perturbative QCD

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Abstract

After some general remarks on non-perturbative QCD I present shortly models which lead to a color-electric flux tube formation. The implications of such a flux tube formation especially on high energy scattering are discussed.

1 Introduction

Few body aspects of non-perturbative QCD cover a wide field, and strictly speaking only meson spectroscopy can be considered to lay outside this domain, since even baryon spectroscopy in non-perturbative QCD is at least a three body problem because of the three valence quarks inside the baryon, not to speak of hadron hadron scattering, which in the simplest case is a four body problem. I can therefore treat only a small subsection of the topics which deserve to be considered. Since furthermore we have no canonical analytical procedure to treat non-perturbative QCD this selection will be necessarily a biased one. My presentation is organized as follows:

1) Some essential features of QCD which are needed for further reference are repeated and the necessity of a non-perturbative treatment in the application to hadron physics is pointed out.

2) A picture which is common to several attempts to understand confinement – the most intriguing feature of non-perturbative QCD – is presented. It is the picture in which a color-electric string is formed between the quarks and (anti-) quarks. Some treatments of non-perturbative QCD which lead to such a string picture are also very shortly exhibited.
3) A simplified version of this picture is given and some consequences of it are discussed, especially baryon spectroscopy and the formation of exotic states.

4) In the last point some typical few body aspects in high energy scattering are presented.

2 Some essential features of QCD

We all tend to agree that the Lagrangian describing strong interactions is known, namely that of quantum chromodynamics, the theory of colored quarks and gluons.

It looks conspicuously similar to that of the time honored theory of electrons and photons, i.e. quantum electrodynamics.

\[
\mathcal{L}_{\text{QCD}}(x) = i \bar{\psi}^{E^f}(x) \gamma^\mu D^C E(x) \psi^{E^f}(x) - m_f \bar{\psi}^{C^f}(x) \psi^{C^f}(x) - \frac{1}{4} F^{\mu\nu F}(x) F^{F}_{\mu\nu}(x). \tag{1}
\]

Here \( \psi^{E^f}(x) \) is the quark field with color \( E \) and flavor \( f \), \( F^{F}_{\mu\nu}(x) \) is the gluon field tensor defined as:

\[
F^{F}_{\mu\nu}(x) := \partial_\mu A^F_\nu - \partial_\nu A^F_\mu + g_s f_{GH} A^G_\mu A^H_\nu \tag{2}
\]

with the color potential \( A^F_\mu(x) \) and the covariant derivative:

\[
D^C E_\mu := \partial_\mu \delta^C E - i g_s / 2 \lambda^C E A^F_\mu \tag{3}
\]

where \( \lambda_F \) are the Gell-Mann matrices.

I just remind you that the nonlinear term in the gluon field-tensor, which is not present in the electro-dynamic field strength tensor, leads through the last term in (1) to cubic and quartic terms in the potentials \( A^F_\mu(x) \) in the QCD-Lagrangian and thus to non-linear terms in the equivalents of the Maxwell equations.

The Lagrangian \[1\] has been constructed in such a way that it is invariant if the quark and gluon fields are transformed by space-time dependent gauge transformations of the group \( SU(3) \). The non-linear terms are a necessary consequence of this local gauge invariance under a non-Abelian transformation.

Analogously to the path integral formulation of quantum mechanics, we obtain the quantum mechanical expectation value of a functional of quark and gluon fields \( F[\psi, A] \) as a functional integral which we write formally as:

\[
< F[\psi, A] > = \frac{1}{\mathcal{N}} \int \mathcal{D}\psi \mathcal{D}\bar{\psi} \mathcal{D}A F[\psi, A] e^{iS_{\text{QCD}}[\psi, \bar{\psi}, A]} \tag{4}
\]

where the normalization constant \( \mathcal{N} \) is given by the integral without the function \( F \):

\[
\mathcal{N} = \int \mathcal{D}\psi \mathcal{D}\bar{\psi} \mathcal{D}A e^{iS_{\text{QCD}}[\psi, \bar{\psi}, A]} \tag{5}
\]
and $S_{QCD}[\psi, \bar{\psi}, A]$ is the action, i.e. the integral of the QCD Lagrangian (1) over the space time continuum:

$$S_{QCD}[\psi, \bar{\psi}, A] = \int d^3xdx_0 \mathcal{L}_{QCD}$$ (6)

If we consider the light quark sector of the theory and go to the chiral limit where all quark masses are zero, the Lagrangian contains no dimensionful quantity and is therefore formally scale invariant. This scale is necessarily broken in order to regularize the Lagrangian and the scale $\Lambda_{QCD}$ is introduced which is the only scale in the chiral limit. The dependence of the coupling constant $\alpha_s \equiv \frac{g^2}{4\pi}$ on an arbitrary energy scale $\mu$ is defined over the so called $\beta$ function:

$$\beta(\alpha_s) = \mu \frac{\partial}{\partial \mu} \alpha_s(\mu).$$ (7)

It can be calculated in perturbation theory and one obtains to lowest order:

$$\beta(\alpha_s) = -\frac{\beta_0}{2\pi} \alpha_s^2 + O(\alpha_s^3) \quad \text{with} \quad \beta_0 = 11 - \frac{2}{3} n_f$$ (8)

where $n_f$ is the number of relevant flavors. The QCD scale $\Lambda_{QCD}$ is now given as an integration constant of the solution of equation (7) and again to lowest order we can write:

$$\alpha_s(\mu) = \frac{4\pi}{\beta_0 \ln(\mu^2/\Lambda_{QCD}^2)} \quad \mu \gg \Lambda_{QCD}$$ (9)

This equation implies asymptotic freedom, since for large scales (high momenta or correspondingly short distances) it vanishes logarithmically.

From this equation we get an intuitive argument that hadronic properties are inaccessible to calculation in perturbation theory. Let us try to calculate for instance a scattering length $a_0$. It can get its dimension only by the the only dimensionful quantity, the regularization scale $\mu$ and hence we would obtain:

$$a_0 = f(\alpha_s) \mu^{-1}$$ (10)

Of course the physical quantity $a_0$ cannot depend on the arbitrary scale $\mu$ and we have:

$$\frac{da_0}{d\mu} = f'(\alpha_s) \frac{1}{\mu} \frac{d\alpha_s}{d\mu} - \frac{1}{\mu^2} f(\alpha_s) = 0$$ (11)

or equivalently:

$$\frac{f'(\alpha_s)}{f(\alpha_s)} = \frac{1}{\beta(\alpha_s)}$$ (12)
which upon integration yields the expression:

\[ f(\alpha_s) = f(\alpha_0) \exp \left( \int_{\alpha_0}^{\alpha_s} \frac{d\alpha'}{\beta(\alpha')} \right) \]  (13)

which cannot be expanded in a power series in \( \alpha_s \) since the powers of \( \alpha_s \) are in the denominator of the exponent.

3 Methods and models in non-perturbative QCD

Unfortunately we have for QCD - as for any other realistic quantum field theory- only one analytical method to treat the functional integral (1), namely perturbation theory. But as mentioned in the previous section, hadronic properties cannot be calculated with perturbation theory, and therefore we have to rely here either on a direct numerical simulation of the functional integral (1) or on models. Models should respect the symmetries and explain the most striking features of QCD: confinement and spontaneous chiral symmetry breaking. The former manifests itself by the fact, that quarks and gluons have never been observed as free particles, the second one, that we have the \( \pi \), \( K \) and \( \eta \) mesons as pseudo-Goldstone particles in the hadronic sector. For the numerical simulations one formulates the classical Lagrangian (1) on a 4-dimensional hyper-cubic lattice in Euclidean space time, preserving gauge invariance (1) and performs the high dimensional integrals numerically. With the appearance of more and more powerful computers this approach gains more and more importance and in some fields, as field theory at finite temperatures it has replaced widely the experimental data. Indeed, lattice people quite often refer to their results as measurements.

The oldest of the models – it is older than QCD – is the non-relativistic potential model of quarks (see e.g. [21]). We can view the process of solving the the Schrödinger-equation of that model as one of the most powerful tools for performing the functional integration. Though potential models give very satisfactory results for hadron spectroscopy, the basic features, confinement and chiral symmetry breaking have to put into the model and are not a consequence of it. More sophisticated methods as attempts to solve the Schwinger Dyson equations or the Bethe Salpeter equation are much closer to the original QCD, but I think it is fair to say that there is not yet a real break-through to note.

Another important attempt to solve the functional integral (1) is based on the expansion of the action around classical solutions, the so called instantons. This approach can explain very nicely chiral symmetry breaking, but not confinement.

An approach based on modifying the QCD equations in order to make them solvable but hopefully preserving the essential features is dual QCD. It goes back to the proposals of t’Hooft [2] and Mandelstam [3] that the confinement mechanism in QCD is structurally similar to the formation of Abrikosov-Nielsen-Olesen strings in a superconductor of the second kind. In the latter case the condensation of Cooper-pairs implies that magnetic fields cannot exist at all in superconductors of the first kind (Meisner-effect) and only as flux tubes of magnetic fields in superconductors of the second kind, see Figure 1a).
Figure 1: The formation of a flux tube in a superconductor of second kind. a) electrodynamics, b)dual QCD

Magnetic monopoles would thus be confined in an electric superconductor of second kind, since the magnetic string between them would lead to a linear increase of the field energy with growing distance. If QCD can be approximated by a dual picture of the electrodynamics of a superconductor of second kind, i.e. if color-magnetic monopoles condense and color-electric flux tubes are formed, color charges are confined, see Figure 1 b). In the classical approximation the equations of motion imply the dual potential $C^\mu$, the monopole condensate $\Phi$ and the field tensor derived from a Dirac string of a quasi-static quark-antiquark pair $G^{S}_{\alpha\beta}$ which yields the coupling of the sources to the fields:

\[
\partial^\alpha (\partial_\alpha C_\beta - \partial_\beta C_\alpha) = j_{\beta}^{\text{mon}} - \partial^\alpha G_{\alpha\beta}^S
\]

\[
(\partial_\alpha - igC_\alpha)^2 = -\frac{200\lambda}{3}\Phi(|\Phi|^2 - B_0^2)
\]

with the monopole current:

\[
j_a^{\text{mon}} = -3ig[\Phi^*(\partial_\alpha - igC_\alpha)\Phi - \Phi(\partial_\alpha + igC_\alpha)\Phi^*]
\]

Monopole condensation (in a maximal Abelian gauge fixing) has been observed on the lattice, thus giving some support to a kind of t’Hooft-Mandelstam mechanism, though confinement in real QCD might be more subtle than in the dual picture \[1\]. There has been a recent revival of interest in dual QCD since in supersymmetric QCD \[2, 3\] the duality transformation can be performed and the long distance behaviour of supersymmetric QCD is described by a weekly coupled dual gauge theory. I cannot go into details of dual QCD and refer to the papers of Ball, Baker, Zachariasen and coworkers (see \[7\] and the literature quoted there) who have elaborated the model in great detail.

The last model I want to present shortly is the model of the stochastic vacuum (MSV) proposed by Yuri Simonov and myself several years ago \[8, 9\], for reviews see \[10, 11\]. It starts from the assumption that the complicated structure of the colordynamic interaction at long distances can be approximated by a simple stochastic process, in the most restricted and phenomenologically most useful form by a Gaussian process, which
is determined by the gauge invariant correlator of two color-field strength tensors \( F^{\mu\nu}_C \) at different space time points.

\[
<F^{\mu\nu}_F(x,y)F^{\kappa\lambda}_F(0,y)> = \frac{\delta^{FG}}{12(N_c^2 - 1)} < FF > \cdot \left\{ (\delta_{\mu\kappa}\delta_{\nu\lambda} - \delta_{\mu\lambda}\delta_{\nu\kappa})D(x^2) \cdot \kappa \right\}.
\]

Here we have switched to an Euclidean space time with \( x^4 = ix^0 \) and \( y \) is a reference point to which the color content has been transported. A nice feature of the model is that it leads to linear confinement for non-Abelian theories, but not to confinement for Abelian gauge theories like QED, unless there is monopole condensation. It also reproduces correctly the spin structure of the confining potential between heavy quarks. There are new developments towards its application to light quarks and an understanding of chiral symmetry breaking \([12]\). Confinement in this model is the result of the special tensor structure in front of \( D(x^2) \) in the correlator \((16)\). This structure is forbidden in an Abelian theory by the homogeneous Maxwell equations, but can be present in a non-Abelian gauge theory like QCD. It had been shown by Di Giacomo and coworkers \([13, 14]\) on the lattice that this confining tensor structure is indeed present and even dominant in QCD.

## 4 Formation of a color-electric flux tube in QCD

String formation between static quarks has been observed in lattice calculations (for SU(2) as gauge group) \([15, 16, 17]\). In Figure 2 taken from reference \([16]\) the action density (i.e. the difference between the squared electric and magnetic field strength) for a static quark antiquark pair at different separations \( R \) is displayed. One clearly sees the contribution of the color-electric Coulomb field around the quark and anti-quark, but also the connecting string is clearly visible.

The formation of a colorelectric flux tube is a direct consequence in dual QCD see Figure 1. Also the model of the stochastic vacuum leads to the formation of a colorelectric flux tube \([17, 18]\). In Figure 3 taken from \([18]\) the squared color-electric field strength as calculated in the model of the stochastic vacuum is displayed. The “Coulomb peaks” around the positions of the quark and anti-quark are due to the Coulomb contribution to correlator \( D_1 \) of equation \((13)\), the connecting string however is entirely due to the typically non-perturbative and non-Abelian structure \( D \).

The excitations of the string, or at least its long distance consequences have been observed in lattice calculations; the two models treat, at least up to now, only a static string. The full width of the string is similar in all cases, about one fermi (a bit less in the SU(2) case). The width of the string is nearly independent of the distance between the quark and antiquark, the same holds for the energy density thus leading to a linearly rising
Figure 2: Action density (in units of the string tension) of a quark antiquark pair at distances of ca 0.7 and 1.35 fm. The figure is taken from [16].

potential energy between the static sources, see Figure 4 for the situation in the model of the stochastic vacuum.

The consequences of string formation for spin and velocity dependent forces between heavy quarks have been studied intensively by Brambilla and Vairo [19] leading to very similar results for the MSV and dual QCD.

If we regard interactions between more than two particles, string formation tells us that we cannot expect that many body forces can be obtained by adding just the two body forces. In general a string is is saturated if it has found its two end points. This implies that in contradistinction to electrodynamics we cannot count with van der Waals forces between hadrons which fall off like an inverse power, at least if the distances between the hadrons are comparable or larger than the string width. This is of course in agreement with the general field theoretical result that forces between hadrons fall off exponentially due to the finite masses of the exchanged hadrons.

But before I come to hadron-hadron interactions let me shortly discuss the nucleon in non-perturbative QCD and exotics.

4.1 Nucleons and exotics

Mathematically a string between colored objects can be expressed by a generalization of the Schwinger string of electrodynamics:

$$S(x, y) := P \exp \left( -\frac{ig}{2} \int_{C(x \rightarrow y)} (A^F_\mu \lambda_F) dx^\mu \right)$$

where $C(x \rightarrow y)$ is a curve connecting the points $x$ and $y$, $\lambda_F$ are for SU(3) the Gell-Mann matrices and $P$ denotes path ordering, that is the prescription how to treat the integral.
over non-commuting quantities in the exponential. $S(x, y)$ is an element of SU(3), and it has the important property to transport the information on color from point $x$ to point $y$. Due to local gauge invariance it makes no sense to speak e.g. of color neutrality of a nonlocal object, like $\sum_C \bar{\psi}^C(x)\psi^C(y)$, unless we insert the Schwinger string $S(x, y)$ between it (or use a gauge, in which $S(x, y)$ is unity).

A special feature of SU(N) is that one can couple $N$ objects transforming under the fundamental transformation to a singlet, in case of SU(3) thus three quarks. This is a trivial consequence of the fact that the determinant of an SU(N) matrix is by definition one, i.e.

$$\epsilon_{iik}U_{ir}U_{ks}U_{lt} = \epsilon_{rst}$$

if $U_{ab}$ is a special unitary $3 \times 3$ matrix. This allows to form a color neutral 3 quark state and was an essential ingredient of the quark respective aces- model of Gell-Mann and Zweig. In the string picture the baryon of three quarks is formed just by joining three Schwinger strings according to (18) at the point $y$:

$$\psi^C(x)\psi^D(u)\psi^E(w)S_{CF}(x, y)S_{DF}(u, y)S_{EF}(u, y)$$

If we assume that the strings carry energy proportional to their length and that the width can be neglected (a rather unrealistic assumption) one can derive the genuine 3-quark potential inside a nucleon, see Figure 3:

$$V_N(\vec{x}, \vec{u}, \vec{w}) = \min_y (|\vec{x} - \vec{y}| + |\vec{u} - \vec{y}| + |\vec{w} - \vec{y}|)$$
Figure 4: a) The width of the string as a function of the distance of the static quark-antiquark pair, the unit $a$ is about 0.3 fm, from [18]. b) The potential energy of a static quark-antiquark pair calculated by the area law of the Wegner-Wilson loop (solid line) and calculated as color-electric field energy (dots) from [18].

The point $y$ for which the rhs of (20) is minimal has been found by Toricelli. If the triangle spanned by the points $\vec{x}, \vec{u}, \vec{w}$ has no angle larger than $2\pi/3$, then $y$ is an inner point of it and its connecting lines to the corners form a star with angles $2\pi/3$. Otherwise, if one of the angles of the triangle is greater than $2\pi/3$ $y$ is the corner point with this angle.

Simple geometrical considerations give

$$
\min_y (|\vec{x} - \vec{y}| + |\vec{u} - \vec{y}| + |\vec{w} - \vec{y}|) = f(\vec{x}, \vec{u}, \vec{w}) \{ |\vec{x} - \vec{u}| + |\vec{u} - \vec{w}| + |\vec{w} - \vec{x}| \} \quad (21)
$$

with

$$
1/2 \leq f(\vec{x}, \vec{u}, \vec{w}) \leq 1/\sqrt{3} (\approx 0.58) \quad (22)
$$

Thus if the string tension is $\sigma$ for the quark antiquark pair, the 3-quark potential in the nucleon (20) can be quite well approximated (to better than 10 %) by the additive effective quark-quark potential

$$
V_N(\vec{x}, \vec{u}, \vec{w}) \approx V_{\text{eff}}(\vec{x}, \vec{u}) + V_{\text{eff}}(\vec{u}, \vec{w}) + V_{\text{eff}}(\vec{w}, \vec{x}), \quad V_{\text{eff}}(\vec{x}, \vec{u}) = 0.54\sigma|\vec{x} - \vec{u}| (23)
$$

These considerations were first put forward by V. Müller and myself [20] inspired from strong coupling expansion of lattice QCD [1]. Calculations of baryonic spectra based on the effective potential (23) turned out to be very reasonable see [21], much more elaborate calculations with very good results were also performed on the basis of similar considerations by Fabre de la Ripelle and Simonov [22].

In the string picture we may also construct more elaborate states than mesons and baryons, as $(qq) - (qq)$ or 6 quark states as indicated in Figure 6.

Spectroscopy has up to now given no compelling reason for the existence of these states. At the moment all existing resonances seem to be compatible to be quarkonia, glue-balls or superpositions thereof [23].
Figure 5: The genuine 3 body potential of the constituent quarks at positions $uvw$ in the baryon as addition of the three string lengths emerging from the point $y$. The dashed lines indicate the approximation by addition of the three lines of the circumference, see equation 21.

Figure 6: Exotic states constructed in the string picture

5 High energy scattering and few-body problems

5.1 Quark-additivity versus string-string interaction

In baryon spectroscopy we have 3 valence quarks involved, in meson baryon scattering at least 5 of them. For low energy scattering one cannot dispose of meson exchange and one comes therefore in problems which are extremely intricate from the QCD point of view. For high energy scattering one expects that the leading contribution to elastic scattering is determined by gluon exchange \([24, 25]\), so there is some hope that with non-perturbative methods one can treat the problem. From the few-body point of view, there is an interesting question: Is the hadron hadron scattering determined by the valence quarks, or does the string also participate at the interaction. In an Abelian Model Landshoff and
Nachtmann [26] found that only the quarks participate and if the distance of the quarks inside a hadron is large as compared to the correlation length of the gluon field strength correlator introduced above, the total hadron cross section is a superposition of the quark cross sections, justifying in such a way the quark additivity found in the early days of the quark model [27, 28]. If we have additivity of the quark cross sections the pion-nucleon cross section is 6 times, the nucleon-nucleon cross section 9 times the quark-quark cross section. The so predicted ratio 2/3 for $\pi^- N$ to $N^- N$ scattering is indeed observed for the total cross sections at high energies.

In the model of the stochastic vacuum, where the non-Abelian nature of QCD plays a crucial role, we can also treat high energy scattering [29], using more or less the same formalism which was applied to find the confining potential or the color-electric flux tube. The essential ingredients are expectation values of two Wegner-Wilson loops with lightlike sides, see Figure 7.

A crucial result of the model is that the same mechanism which leads to string formation and confinement also leads to a string-string interaction, thus invalidating quark-additivity. Nevertheless also this model yields the correct ratio of $\pi^- N$ to $N^- N$ scattering.
scattering without any new free parameter, since it is determined by the known (electromagnetic) radii of the pion and the nucleon. It furthermore predicts the ratio of $K - N$ to $\pi - N$ scattering correctly also from the electromagnetic radii, without introducing a different cross section for scattering of strange and non-strange quarks as has to be done in the quark additivity scheme and which is strange to the flavor independence of the quark-gluon interaction, see Table 1. An essential feature of the string-string interaction is the increase of the cross section with increasing hadron size. In electro-production of vector-mesons the virtuality $Q^2$ of the photon is the handle for constructing quark-antiquark states with different sizes $R_{q\bar{q}} \propto \frac{1}{Q}$. Nemchik et al. [30] have disentangled the cross section of a quark-antiquark state (color-dipole) of a certain size $R_D$ from data of electro-production of vector mesons from protons. Their results are shown in Figure 8. The dashed line is their phenomenological fit to the dipole-proton cross section, the solid line is our prediction from the model of the stochastic vacuum, where no high-energy data have been used [31]. The data show no sign of saturation at dipole sizes small as compared to the hadron radii, as would be necessary for quark additivity to hold and our prediction fits the semi-experimental results within the expected accuracy.

### 5.2 Nucleon Structure and the C=P=-1 Exchange

If the leading contribution to high energy scattering is due to gluon exchange, there is no compelling reason, why the exchange with vacuum quantum numbers, i.e. the one with positive parity and positive charge parity (C=P=+1) should be overwhelmingly dominant. Indeed if the leading contribution to the C=P=+1 exchange is a “non-perturbative two gluon state” a “non-perturbative three gluon state” should not be terribly suppressed and such a state can also have the quantum numbers C=P = - 1. The exchange of such an object has been considered first by Lukaszuk and Nicolescu [32] and been called odderon. Such a state would contribute differently to proton-proton and proton-antiproton scattering and lead to a difference in the ratio of the real to the imaginary part of the respective scattering amplitudes. However such a difference has not been observed, experimentally the difference

$$\Delta \rho(s) = \frac{\text{Re}[T^{pp}(s, t = 0)]}{\text{Im}[T^{pp}(s, t = 0)]} - \frac{\text{Re}[T^{pp}(s, t = 0)]}{\text{Im}[T^{pp}(s, t = 0)]}$$

is at $\sqrt{s} = 546$ GeV [33] smaller than approximately 0.02 and well compatible with zero. This indicates a strong suppression of the odderon. Such a behavior is hardly compatible with an additive quark model, and indeed Donnachie and Landshoff [33] have found in an Abelian model for the pomeron [20] a value of $\Delta \rho$ of about 0.5. The model of

| Ratio                  | Model | Exp. |
|------------------------|-------|------|
| $\sigma_{\pi N}/\sigma_{pp}$ | 0.66 ± 0.02 | 0.63 |
| $\sigma_{pK}/\sigma_{p\pi}$ | 0.82 ± 0.08 | 0.87 |

Table 1: Predictions of the model of the stochastic vacuum for total $\pi$-nucleon and $K$-nucleon cross sections, from [29]
Figure 8: Comparison of our result for the total cross section for dipole (extension $R_D$) proton scattering with values extracted from cross sections of lepto-production of vector-mesons by the method of Nemchik et al. [30]. The solid line is our result without any fitting of the parameters to high-energy data. The dashed line is the ansatz of Nemchik et al. for the total cross section, [31].

the stochastic vacuum applied to high energy scattering would yield a similar value if the nucleon is assumed to be a fully symmetric 3 quark structure. There are however several indications that two quarks in the nucleon cluster to a diquark. In such a case this diquark behaves as far as color is concerned like an antiquark. The odderon may couple to such a state, but if we rotate the quark-diquark state by $\pi$ we obtain, due to the negative parity of the odderon, a contribution of opposite sign which cancels the unrotated one. The effect has been discussed quantitatively in reference [34] and the following results have been found. If the three quarks sit at the corners of an triangle and we decrease the baseline $r_{\perp}$ the $C=P=-1$ contribution decreases dramatically leading to a $\Delta \rho \leq .02$ for a diquark radius of ca. 0.3 fm, see Figure 9. Thus even a moderately extended diquark yields the desired odderon suppression. This mechanism can be tested further at Hera. If one considers the photo or electro-production of pseudo-scalars and pseudo-vectors the odderon will couple to the photon-meson transition. If the nucleon has the diquark structure as discussed above, it cannot couple to a proton directly, but to a proton-baryon($J^P = \frac{1}{2}^{+}$) transition vertex, see Figure 10. A good candidate for the isolation of the odderon would thus be a reaction like:

$$\gamma^{(*)} + p^{-} \rightarrow \eta + N(1535)(\frac{1}{2}^{-})$$

in the diffractive region. The signatures would be the decay gammas of the $\eta$ and the decay neutron of the N(1535).
Figure 9: The difference $\Delta \rho$, see equ. [24] as a function of the “diquark radius” $r_\perp$, [34].

Figure 10: Odderon exchange contribution to diffractive photo- or electro-production of odd parity states off nucleons

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