PROCEEDINGS

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High energy scattering and vacuum structure

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ABSTRACT: The implications of typically nonperturbative features of QCD on high energy reactions are discussed. Special emphasis is given to the role of a colour-electric string in scattering processes.

KEYWORDS: QCD, nonperturbative effects, confinement.

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1. General remarks

The novel feature in particle physics which emerged in the second half of the twentieth century is undoubtedly confinement. It still remains an unsolved problem, but there is strong evidence that the physical reason for confinement is explained by the time-honored t'Hooft-Mandelstam picture [1, 2].

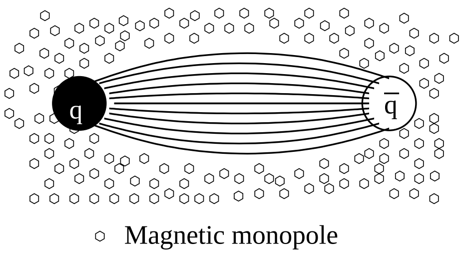

Figure 1: Monopole condensation and colour-electric flux tube formation in the QCD vacuum

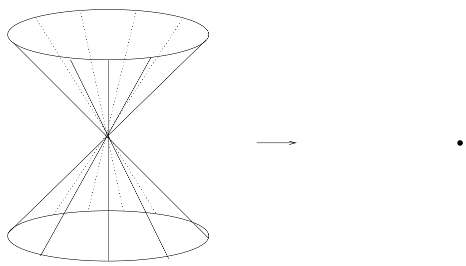

Figure 2: The fate of the light cone upon transition from Minkowskian to Euclidean metric

*Speaker.

Monopoles condense to form the physical vacuum and force the colour-electric field lines into a flux tube. This is analogous to the formation of magnetic flux tubes in a type II superconductor due to condensation of Cooper pairs. It leads to the interesting question: Which influence has the presumably non-trivial structure of the vacuum and especially the flux tube on high energy scattering of hadrons.

I first shall give a very sketchy overview of what we know about the vacuum. The only answer to this question which is based on first principles, that is the QCD Lagrangian, comes from numerical simulations of the lattice regularized version of QCD. Many investigations have been devoted to this problem, especially by Di Giacomo and collaborators from the Pisa Group [3]. All are consistent with the t'Hooft-Mandelstam picture. It is however an extremely difficult task to get information on high energy scattering from lattice QCD and this has a very simple kinematical reason. High energy scattering is governed by dynamics near the light cone $\vec{x}^{\,2} - c^2t^2 = 0$. Lattice QCD has to be performed in Euclidean metric with the substitution $t \to x_4 = ict$ and hence the light cone shrinks to the point $\vec{x}^{\,2} + x_4^2 = 0$, see Figure 2. Therefore one has to continue from a tiny region to the full range of scattering dynamics. This will, if possible at all, be a very difficult task. This all will be treated in a devoted talk at this conference by E. Meggiolaro [4]. A more immediate way to get from lattice calculations information related to high energy scattering is the investigation of glueball spectra, which should lie on Regge trajectories. This will be treated in the talk of M. Teper [5].

A rather save knowledge we have of the vacuum is that there exist instantons. These are minima of the classical non-Abelian Lagrangian in Euclidean geometry. These solutions are inherently nonperturbative, they are of the form:

$$\mathbf{A}_\mu = \frac{x^2}{x^2+\rho^2}\partial_\mu g(x) g^{-1}(x) \quad g(x) = \frac{x_0 + i\vec{\sigma}.\vec{x}}{\sqrt{x^2}} \tag{1.1}$$

They are typical for the gauge group $SU(2)$ as can be seen from the occurrence of the Pauli matrices σ but also exist for the QCD Lagrangian since $SU(2)$ is a subgroup of $SU(3)$. The size of an instanton is determined by the size parameter ρ and a crucial function is the instanton size distribution $D(\rho)$. For small size instantons the distribution function can be calculated from first principles and the contribution to high energy scattering has been intensively discussed by Ringwald and Schrempp [6, 8, 10, 9]. A dedicated H1-experiment [11] at DESY finds that instantons improve the fit compared to conventional Monte-Carlo simulations which, however does not prove definitely their contribution to high energy reactions.

Large size instantons have to be treated by models, as the instanton liquid model of Shuryak. Since instantons lead to high multiplicities in particle production they might be quite essential for the increase of cross sections with high energy [12, 13, 14], this will be treated in the contribution of E. Shuryak to this conference [7].

Another source of information we have about the QCD vacuum comes from the sum-rule technique developed by Shifman, Vainstein and Zakharov [15]. It is based on the operator product expansion in quantum field theory. It is proven to all orders in perturbation theory and assumed to be valid beyond it. It leads to power corrections to purely perturbative results of vacuum expectation values of products of operators

$$\langle A_1(x)A_2(y)\rangle = \langle : A_1(x)A_2(y) :\rangle_{\rm pert} + \sum_i F_i\Big((x-y)^2\Big)\langle O_i(x)\rangle \tag{1.2}$$

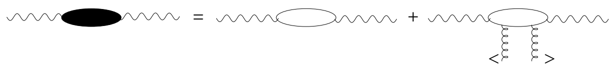

Figure 3: Contribution of the gluon condensate to the vacuum expectation value of a current of a quark and an antiquark

The possible structure of the local operators $O_i(x)$ is determined by the operator products and the coefficient functions $F_i\Big((x-y)^2\Big)$ can be calculated in perturbation theory, but the values of the condensates, $\langle O_i(x)\rangle$, are phenomenological input describing certain aspects of the vacuum structure. If the operator $O(x)$ is the gauge and Lorentz-invariant normal ordered product $: g^2/2 \text{ tr } \mathbf{F}_{\mu\nu}\mathbf{F}^{\mu\nu} :$, the contribution of the condensate to the vacuum expectation value of a current can be visualized by the picture given in Figure 3

An inherent difficulty of this approach is that the perturbative part, indicated above by the normal ordering :: cannot be defined clearly. A certain pragmatism is necessary and has let to impressive results. From many analyses the vacuum expectation value of the above mentioned gluon operator, the gluon condensate, has been determined to, see e.g. [16]

$$\langle : g^2/2 \text{ tr } \mathbf{F}_{\mu\nu}\mathbf{F}^{\mu\nu} : \rangle = (1 \pm 0.15 \text{ GeV})^4 \tag{1.3}$$

Here and in the following **bold** letters denote matrix valued quantities for instance $\mathbf{F}_{\mu\nu} = \sum_c F^c_{\mu\nu}\lambda^c/2$.

The presence of a colour-magnetic field in the vacuum leads to interesting effects as has been stressed by O. Nachtmann and collaborators [17, 18]. In the perturbative treatment of virtual photons through the Drell-Yan process factorisation is assumed, but since the quark and the antiquark travel through the same magnetic field (Figure 4) their spins get correlated and violate factorisation. Such a violation of the standard factorisation ansatz has been indeed observed in the reaction $\pi^- p \to X\gamma^*_{\hookrightarrow \mu^+\mu^-}$

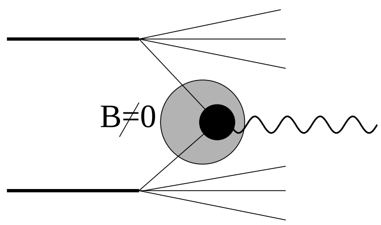

Figure 4: Violation of factorisation through partial spin alignment due to the presence of a colour-magnetic vacuum field

Neither instantons nor condensates can explain confinement. In order to investigate deeper a possible influence of the confinement mechanism on high energy reactions one has to rely on models. In the following I shall present an approach to soft high energy scattering developed in Heidelberg by Nachtmann, Pirner, myself and many excellent graduate students. In order to investigate nonperturbative effects on scattering one has to do two things: First to develop a formalism which allows to calculate nonpertubatively amplitudes of high energy reactions, at least in principle and second to perform then the actual calculation. The first step has been done by Nachtmann in 1991 and I shall very shortly outline the principle ideas, for details I refer to the original literature [19] and to reviews [22, 23, 24, 21].

2. Nonperturbative scattering

The formalism of Nachtmann is based on the functional integral approach to quantum field theory. According to it the vacuum expectation value of a function of a quantized field $\mathbf{A}$ is expressed formally as the functional integral:

$$\langle f(\mathbf{A})\rangle = \int \mathcal{D}\mathbf{A} f(\mathbf{A}) \exp(-iS_{QCD}[\mathbf{A}]) \tag{2.1}$$

In order to treat scattering of two highly energetic quarks one first considers the scattering of them in an external field. Making use of the WKB approximation the scattering amplitude is to leading order in energy given by the phases $\prod_{i=1,2} e^{-ig\int_{P_i}\mathbf{A}dx}$ picked up by the quarks along their classical lightlike paths $P_i,\ i=1,2$, see Figure 5

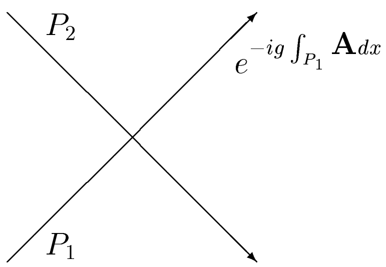

Figure 5: Scattering of a highly energetic quark and an antiquark in WKB approximation

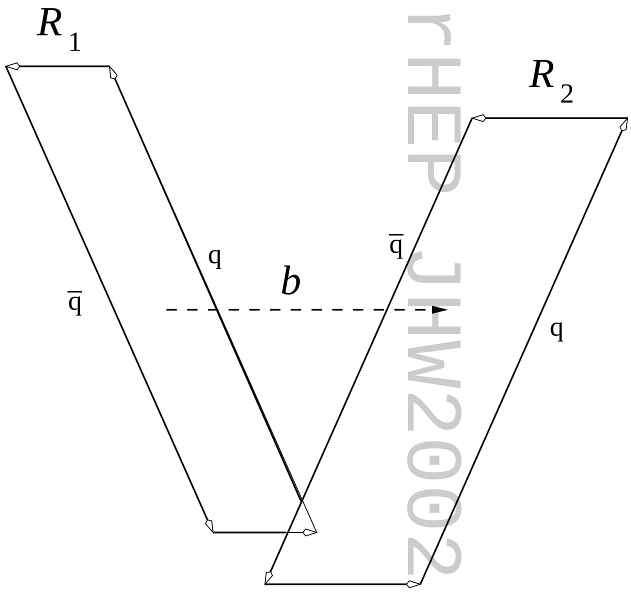

Figure 6: Scattering of two dipoles with dipole vectors R_i and impact parameter $\vec{b}$

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The amplitude for scattering in a quantized field is then given by the functional integral over these phases with the exponential of the action as weight.

$$\langle e^{-ig\int_{P_1}\mathbf{A}dx} e^{-ig\int_{P_2}\mathbf{A}dx}\rangle = \int \mathcal{D}\mathbf{A} \prod_{i=1,2} e^{-ig\int_{P_i}\mathbf{A}dx} \exp(-iS_{QCD}[\mathbf{A}]) \tag{2.2}$$

This expression is not gauge invariant and therefore one has to fix the gauge to make it well defined. If we are however interested in the effects of confinement it is essential not to consider the scattering of colour siglets and not of hypothetical free quarks. It was therefore proposed [20] to consider Wegner-Wilson loops with light-like sides rather than Wilson lines as building blocks for high energy scattering. This leads to the scattering amplitude of two dipoles with vectors $\vec{R}_1,\ \vec{R}_2$ and impact parameter $\vec{b}$:

$$J(\vec{b},\vec{R}_1,\vec{R}_2,z_1,z_2) = \langle \mathrm{Tr}\ e^{-ig\int_{W_1}\mathbf{A}dx}\mathrm{Tr}\ e^{-ig\int_{W_2}\mathbf{A}dx}\rangle \tag{2.3}$$

visualized in Figure 6. A loop is composed of the light-like paths of the quark and the antiquark inside a hadron and the strings connecting them in the remote past and future.

Unfortunately even the definition of the functional integral is only well understood for Gaussian measures, that is an action S which is quadratic in the stochastic variable $\mathbf{A}$. There are essentially two a priori methods to treat more general processes:

- Numerical evaluation of the integrals in the lattice regularized version.
- Perturbation theory: The action is split into a quadratic part and the rest

$$S_{QCD}[\mathbf{A}] = S_0[\mathbf{A}] + ig_s S_{\mathrm{int}}[\mathbf{A}] \tag{2.4}$$

and the exponential measure $\exp(-iS[\mathbf{A}])$ is expanded in g_s,

$$\exp(-iS_0[\mathbf{A}]) = \exp(S_0[\mathbf{A}])\Big(1 - ig_s S_{\mathrm{int}}[\mathbf{A}] + \dots\Big) \tag{2.5}$$

Now the functional integrals are Gaussian and can be performed. This method was used originally by Feynman in order to derive the perturbative series and led him to his famous rules and diagrams.

Unfortunately both methods seem not to be very useful for our purpose, the lattice calculation because of the difficulties to treat high energy reactions on an Euclidean lattice as mentioned above, perturbation theory since we are interested in nonperturbative effects.

3. The stochastic vacuum model and high energy reactions

We therefore have to rely on models and in the following I shall concentrate on an approach based on a model proposed by Yuri Simonov and myself [25, 26]. I emphasize that the version treated here needs more assumptions than the original version and I shall call it the extended version of the stochastic vacuum model. Again I give only a very short outline and refer to the literature [22, 23, 24, 28, 21].

The basic assumption of the model is: The long range part of QCD can be approximated by a **Gaussian** stochastic process in the **Gluon Field Strength** $\mathbf{F}_{\mu\nu}$

The implications are most easily seen in the cumulant (linked cluster) expansion of expectation values:

$$\langle \text{ Tr } e^{-ig\int_{\mathcal{F}} d\sigma_{\mu\nu}\mathbf{F}_{\mu\nu}(x)}\rangle \;=$$
$$\exp\left[-\frac{1}{2}g^2\int_{\mathcal{F}}\int_{\mathcal{F}} d\sigma_{\mu\nu}d\sigma'_{\kappa\lambda}\langle\, \mathbf{F}_{\mu\nu}(x)\mathbf{S}(x,x')\mathbf{F}_{\kappa\lambda}(x')\rangle \;+ \text{higher cumulants}\right] \quad (3.1)$$

The path ordered generalized Schwinger string $\mathbf{S}(x,x') = \exp\Big(-ig\int_x^{x'} dz.\mathbf{A}\Big)$ has to be introduced to ensure gauge invariance.

The Gaussian assumption corresponds to the assumption that all higher cumulants in 3.1 are zero. Therefore, in the model the long range part of QCD is determined by the nonlocal correlator of two gluon fields. All vacuum expectation values of gauge invariant expressions factorise into products of these correlators. The correlator is charaterized essentially by two quantities, its strength given by the gluon condensate and its correlation length.

The model is especially convenient to evaluate Wegner-Wilson loops, since the line integrals can be transformed into surface integrals using the non-Abelian Stokes theorem

$$\langle \text{Tr } e^{-ig\int_C \mathbf{A}_\mu(x)\, dx_\mu}\rangle \overset{n.A.Stokes}{=} \quad (3.2)$$

$$\langle \text{Tr } e^{-ig\int_W \mathbf{F}_{\mu\nu}(x)\, d\sigma_{\mu\nu}}\rangle \overset{Gaussian}{=} e^{-\frac{g^2}{2}\int_W\int_W\langle\mathbf{F}_{\mu\nu}(x)\,\mathbf{S}(x,x')\mathbf{F}_{\kappa\lambda}(x')\rangle d\sigma_{\mu\nu}\, d\sigma'_{\kappa\lambda}} \quad (3.3)$$

There are several nice features of the SVM:

• It yields **confinement** for **non**-Abelian gauge theories; for Abelian theories only if magnetic monopoles condense. It also yields a very reasonable value for the string tension, $\sigma = 0.17$ GeV^2 if the correlator of obtained from lattice QCD [29] is inserted.

• The resulting dynamical picture is string formation [30]in conformity with the t'Hooft-Mandelstam picture. The model is consistent with low energy theorems [34].

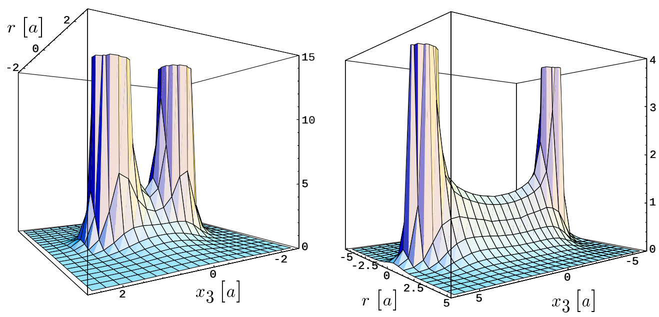

Figure 7: The field density of a static quark-antiquark pair, on the left figure the separation is 2 and on the right figure 9 correlation lengths, from [30]

•It yields Casimir scaling for loops in different representations of $SU(3)$, confirmed by lattice results [31, 32].

There is however also a very **unpleasant** feature: a Gaussian a process for **non-commuting** variables is not uniquely defined.

We now apply the model to high energy scattering by evaluating the expectation values of the Wegner Wilson loops in 2.3 using the model [20, 33]. Though the model has been formulated in Euclidean metric, it is easily continued to Minkowski space [23]. This has the following reason: If we Fourier-transform the correlator the integration over the light-like paths in Figure 6 leads to δ-distributions $\delta(k_0 + k_1)$ and $\delta(k_0 - k_1)$, hence only the transverse components k_2, k_3 survive. These components are not affected by continuation from Minkowskian to Euclidean metric or vice versa and therefore the Euclidean results for correlators (from lattice e.g.) can be used for high energy scattering. It has been shown recently that also a continuation of tilted loops in Euclidean space leads to the same result [34].

Another nice feature is that for a single loop with light-like sides one obtains:

$$\frac{1}{N_c}\langle \mathrm{Tr}\; e^{-ig\int_{W_1}\mathbf{A}dx}\rangle = 1 \tag{3.4}$$

and therefore we have no problems with the quantum field theoretical wave function renormalization.

We next come to the question how to treat the two loops ?

• There is a pedestrian way: one expands the exponentials and applies Gaussian factorisation [33].

• A more refined method was developed in [36]. The product of **two** traces, each in $SU(3)$ is expressed in **one** TRACE in $SU(3) \otimes SU(3)$, formally

$$\mathrm{Tr}\; e^{-ig\int_{W_1}\mathbf{A}dx}\mathrm{Tr}\; e^{-ig\int_{W_2}\mathbf{A}dx} = \mathrm{TR}\; e^{-ig\int_{W_1}\mathbf{A}dx \oplus (-ig)\int_{W_2}\mathbf{A}dx} \tag{3.5}$$

and factorisation of the stochastic variables valued in $SU(3) \otimes SU(3)$ is assumed. This procedure guarantees unitarity. Both methods need an extension of the original version of the model. If one wants to calculate one loop one has to assume factorisation of **matrices** $\mathbf{F}_{\mu\nu}$. For two or more loops one must assume factorisation of the **colour components** $F^c_{\mu\nu}$. Furthermore, the application of the non-Abelian Stokes theorem to two loops is far from trivial. Here special choices for the resulting surface have to be made. This choices influence the numerical results and lead to different sets of input parameters but do not affect the general results.

3.1 General results

We consider first the dipole-dipole cross section

$$\sigma(R_1, R_2) = \int d^2b \int \frac{d\phi_1}{2\pi}\frac{d\phi_2}{2\pi} J(\vec{b}, \vec{R}_1, \vec{R}_2) \tag{3.6}$$

The same mechanism which leads to confinement also leads to **string-string interaction in high energy scattering** which manifests itself in a linear increase of the cross section with dipole size as can be seen in Figure 8. It should be noted however that internal fermion loops will lead eventually to string breaking and the cross section will level off, like the confining potential. This is expected to happen between one and two fm.

The order of magnitude comes out reasonable if both radii are put to typical hadronic sizes of somewhat below 1 fm without any recurrence to high energy parameters.

The resulting dipole-dipole cross section (3.6) factorises to good approximation and a convenient numerical parametrization is

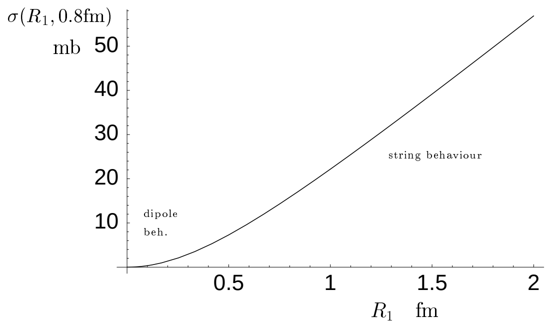

Figure 8: The dipole-dipole cross section for one dipole fixed at $R_2 = 0.8$ fm

$$\sigma(R_1, R_2) = 0.67\frac{1}{4\pi^2}(< g^2FF > a^4)^2 \times R_1\left(1 - e^{-\frac{R_1}{3.1a}}\right)\Big) R_2\left(1 - e^{-\frac{R_2}{3.1a}}\right)\Big)$$

The parameters fine-tuned for p-p scattering are:

correlationlength $a = 0.35\text{fm} \quad \langle g^2FF\rangle a^4 = 23.8$

If one dipole is very small, QCD factorisation should hold: All nonperturbative effects should be accounted for by properties of the larger dipole. But even for a very small dipole scattering off a large one the string contribution is still important. This looks like a violation of factorisation but it is not. It has been shown recently [35] that also in the model the string part of the interaction can indeed be pushed into the nonperturbative part of the large dipole using only mathematical identities. If both dipoles are small, the string contribution in the model is negligible.

3.2 Comparison with experiment

In order to apply the model to hadron-hadron scattering one forms a superposition of dipole-dipole amplitudes with light-cone wave functions of hadrons as weights. For a reaction 1 2 → 3 4 one obtains the expression:

$$T(s,t) = 2isintd^2be^{i\vec{q}\cdot\vec{b}} \int d^2R_1dz_1 \int d^2R_2dz_2\ \psi_3^*(R_1, z_1) \times \psi_1(R_1, z_1)\psi_4^*(R_2, z_2)\psi_2(R_2, z_2)J(s, \vec{b}, \vec{R}_1, \vec{R}_2, z_1, z_2)$$

For hadron wave functions Gaussians were used which are normalised to one – this is necessary for consistency – and the width is adjusted to measurable quantities as the electromagnetic radius or decay constants. For baryons a three quark structure [33] could be used but most applications have been made so far in the much simpler quark-diquark picture. For real or virtual photons wave functions [37] of the form as obtained in perturbation theory are used, but with a quark mass $m(q^2)$ which was adjusted to the two point function of the corresponding interpolating field operator. In model investigations it has been shown that this simple procedure yields indeed a reasonable description of confinement effects in the spacelike region. The model has been applied to many processes and I show here only a few examples.

In Figure 9 the Pomeron contribution to total hadron-hadron cross sections at $\sqrt{s} \approx 20$ GeV is shown. The famous ratio $\sigma_{\pi N} : \sigma_{NN} \approx 2:3$ comes out correctly in the model but not as a consequence of the quark counting rule which is not valid in the model, but as a consequence of the different sizes of mesons and nucleons. Also the ratio $\sigma_{KN} : \sigma_{\pi N} \approx 0.85$ comes out correctly in the model.

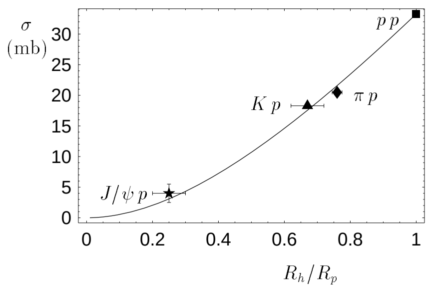

Figure 9: Pomeron contribution to total hadron-hadron cross sections at $\sqrt{s} \approx 20$ GeV from [21]

Since the model leads to a representation in impact parameter space, one can easily obtain differential cross sections. For differential cross sections for $|t| > 0.5$ GeV^{2}the matrix cumulant method [36] is essential. In Figure 10 I show results obtained in [38].

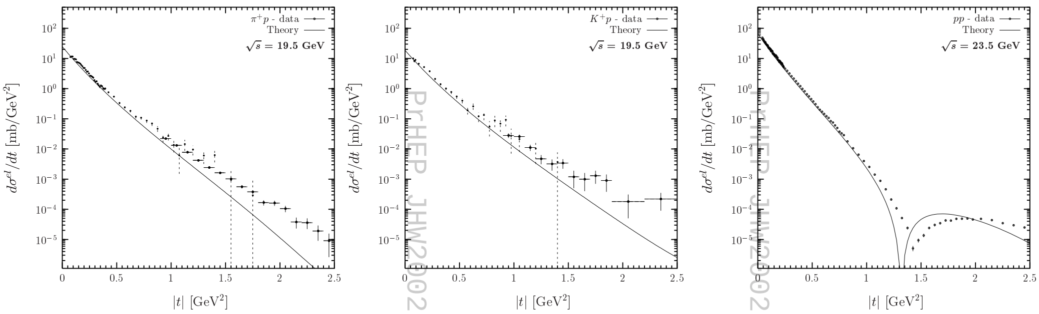

Figure 10: Differential cross sections for πp, Kp and pp elastic scattering at $\sqrt{s} \approx 20$ GeV, from [38]

The model as it stands yields no energy dependence, it has to be put in by hand. One way is to adopt the two Pomeron model of Donnachie and Landshoff [39] and to couple small dipoles ($R \leq R_C \approx 0.2$fm) to the hard Pomeron with intercept 0.42 and large dipoles ($R > R_C$) to the soft Pomeron with intercept 0.08. In this way many phenomenological features can be explained: The peculiar Q^2 dependence of the residues of the soft and the hard pomeron, the hard behaviour of the charm contribution to the proton structure function, the difference between this contribution and the softer energy dependence of J/ψ photoproduction among others. As an example I show the photon-proton cross section for different values of the photon virtuality Q^2 over the measured energy range at HERA.

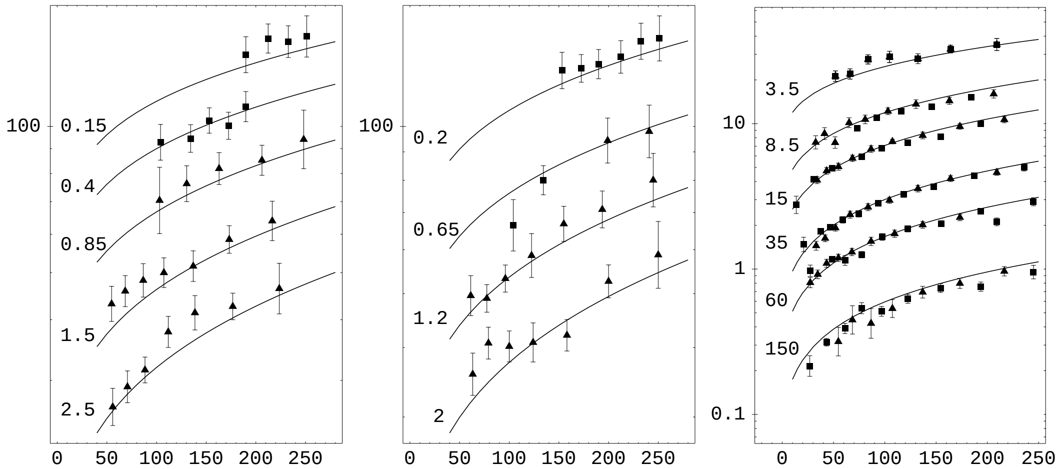

Figure 11: $\gamma^* p$ cross sections $\sigma_{\gamma^* p}/\mu b$ for different photon virtualities Q^2/GeV^2 as function of the energy W/GeV, from [40]

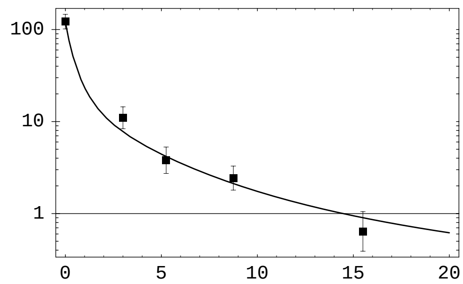

Figure 12: Cross section for the reaction $\gamma^* p \to \gamma p$ in nb at $< W >= 75$ GeV as function of the virtuality Q^2/GeV^2 of the incoming photon, from [41]

In Figure 12 the virtual Compton scattering cross section $\gamma^* p \to \gamma p$ is shown and compared with experiment (after subtraction of the purely electrodynamical contribution). We see also in that case where ,,generalized gluon" distributions occur, a very satisfactory agreement with experiment from the nonperturbative to the perturbative region.

In general the model gives a unified description of scattering amplitudes and has been applied to: pp, πp, Kp, $J/\psi p$ total, differential elastic and diffractive cross sections, $\gamma^{(*)}p$ reactions, photo- and electroproduction, $\gamma^{(*)}\gamma^{(*)}$ reactions. It has also been used to look into the microscopic structure of scattering, for instance by calculating unintegrated gluon distributions [35].

Generally agreement with experiment is on the 20 % level or better. The charm contribution to the photon structure function however comes out by nearly a factor 2 smaller than present experiments. A serious problem is the absence of any sign of odderon contributions in photoproduction of pseudoscalar and tensor mesons. In this case the exchanged trajectory for diffractive scattering must have C-parity -1, that is it must have Odderon quantum numbers. From the point of view of QCD such a trajectory occurs naturally since the symmetric coupling of three gluons just gives such a state with C-parity -1. The weak or absent Odderon contribution to pp scattering can be explained by diquark clustering, but no such excuse is possible for diffractive meson production. If the proton can break up into a state with opposite parity there is no special suppression mechanism at work even in the quark-diquark picture. Cross sections for photoproduction of the π^0, $f_2(1270)$ and the $a_2^0(1320)$ have been calculated in the model [42, 43]. They are displayed in Table 1, third row. The numbers were big enough to motivate an experimental search at HERA, the experiment [44] will be described in the contribution by Kiesling. I only quote the numbers in Table 1, second row. There is a huge discrepancy between prediction and experiment especially for the π, it becomes even more marked if one looks at the spectra.

The discrepancy is astonishing, since an Odderon contribution is expected from QCD. Possible reasons for the failure are:

• The $\gamma\pi^0 O$ coupling is overestimated in the model and the Goldstone boson nature of the pion lets it play a special role. This argument does not apply for a_2, f_2 production since these are „normal" mesons. Therefore better data for f_2 and a_2 production are important.

Reaction	**H1**	Theor.
$\gamma p \to \pi^0 N^*$	$<$ **38**	294 ± 150
$\gamma p \to f_2 N^*$	$<$ **12**	21 ± 10
$\gamma p \to a_2^0 N^*$	$<$ **62**	190 ± 100

Table 1: Experimental limits (95% c.l.) from H1 [44] and theoretical predictions [42, 43] for photoproduction with Odderon exchange

• The Odderon and Pomeron intercepts come out to be one in the model. There are however investigations [45] based on the glueball spectra which predict an intercept of the $C = -1$ trajectory of -1.5, which would make it completely unobservable. Indeed, if one takes the glueball masses from lattice calculations and assumes – for reason of no better knowledge – that the Odderon trajectory is parallel to the Pomeron trajectory with a slope of 0.25 GeV^{-2}, one obtains an Odderon intercept near -1, as can be seen in Figure 13. The Pomeron intercept comes out to small, but as can be seen the Odderon intercept is much lower. In order to explain the absence of a signal in the H1 data of π^0 production an intercept of the Odderon $\alpha_O(0) \leq 0.5$ is sufficient.

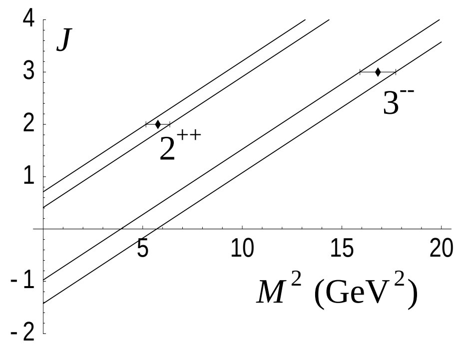

Figure 13: Lattice results for glueball masses on the Pomeron and odderon trajectory. The lines have a slope of 0.25 GeV^{-2}

4. Conclusions

We have seen that the structure of the QCD vacuum can have many effects on high energy reactions. Calculations based on first principles are difficult model investigations are encouraging. Small size instanton effects are model independent and at the threshold of observability, large size instanton effects are model dependent but may explain important features of high energy scattering as the soft energy dependence in hadronic interactions. The presence of a magnetic field in the vacuum, due to gluon condensation, leads to an violation of factorisation in Drell-Yan processes.

In the investigations based on the stochastic vacuum model it turns out that low energy parameters as determined from spectroscopy or lattice calculations determine an important part of high energy reactions. Furthermore it turns out that the same mechanism which leads to confinement also leads to an important contribution to high energy reactions via string-string interaction. In general the agreement between the results of the model and experiment is very satisfactory for a very broad range of reactions but there is a serious discrepancy for diffractive processes where the odderon is exchanged.

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