Vol. 34 (2003) ACTA PHYSICA POLONICA B No 6

SOFT POMERON TRAJECTORY FROM A SATURATION MODEL

Leszek Motyka

M. Smoluchowski Institute of Physics, Jagellonian University
Reymonta 4, 30-059 Kraków, Poland

(Received May 5, 2003)

Dedicated to Jan Kwieciński in honour of his 65th birthday

A saturation model of high energy scattering is analysed, taking into account the impact parameter dependence. The unitarisation is assumed to occur independently at each transverse location. In the single scattering regime the amplitude has features characteristic for the hard Pomeron, and at very high energies the soft Pomeron trajectory is obtained. It is shown, how the string tension may be embodied into the saturation model, linking the Regge phenomenology, the hadron structure and the saturation model.

PACS numbers: 13.85.Dz, 13.85.Lg

1. Introduction

Accurate description of hadron scattering at high energies continues being a challenge. Only the perturbative regime, probed *e.g.* in deep inelastic scattering experiments, is well understood. Amplitudes for soft processes are still incomputable from first principles. Thus, much activity has been recently directed towards extending the perturbative results for semihard processes into the non-perturbative domain. A reasonable extrapolation of this kind must be based on general features of the theory. For the high energy scattering, the guideline has been provided by the unitarity constraints.

In the perturbative picture of hadronic interactions the total cross section is dominated by the exchange QCD Pomeron. In the leading logarithmic approximation the imaginary part of the forward scattering amplitude is given by the BFKL equation [1]. Solutions to this equations exhibit a power like dependence on the collision energy, s^{α_P}, with the Pomeron intercept $\alpha_P \sim 1.3$. Consequently, the increase of the total cross section with energy, $\sigma \sim s^{\alpha_P - 1}$, if continued, would eventually lead to a violation of the Froissart unitarity bound [2] valid for theories with no massless excitations.

The apparent violation of unitarity is a consequence of the approximation, restricted to a single exchange of the BFKL pomeron. Beyond this approximation, multi-pomeron exchange and multi-pomeron interaction should be taken into account. These contributions increase with energy faster than the hard pomeron amplitude and a resummation is necessary. The resummation of leading terms was performed in perturbative QCD by Balitsky [3] and Kovchegov [4] in a special case of a small onium scattering off a large target (nucleus). Solutions to the Balitsky–Kovchegov (BK) evolution equations give unitary amplitudes.

It was assumed in most studies of the BK equation that the dependence on the impact parameter may be neglected. In this case the equation gives a unitary S-matrix, with no diffusion in the impact parameter [4–8]. It had been overlooked until recently, however, that the correct treatment of the BK equations implies a power-like growth of the interaction region with energy [9]. Clearly, this result contradicts the Froissart bound, in spite of unitarity of the scattering matrix at each impact parameter. The conflict is caused by the assumption that the gluon range is infinite, used in the derivation of the BK kernel. The discrepancy disappears when the confining properties of the QCD vacuum are taken into account, restricting the gluon range to distances smaller than one fermi. Thus, the parameters of confining vacuum necessarily enter the high energy scattering amplitudes.

Analysis of the BK equation inspired the Golec-Biernat–Wüsthoff (GBW) saturation model [10], which is a surprisingly simple and efficient parametrisation of the colour dipole cross section, assuming unitarisation of the hard Pomeron. With this model, low-x data from HERA on the total virtual photon–proton cross section were successfully fitted for all virtualities down to the photoproduction limit. Even more surprisingly, the diffractive cross section was also correctly reproduced. In addition, the two-photon cross sections at high energies were fitted by a generalised saturation model [11]. We may conclude that the saturation model provides a sensible treatment of both the *perturbative* and the *nonperturbative* regime of high energy scattering. A natural question arises, whether the saturation model may give insight into a deeply non-perturbative phenomenon of the soft Pomeron exchange, see *e.g.* [12, 13].

When reconciling the saturation model with the Regge model one encounters a subtle problem. Namely, linearity of Regge trajectories has a natural explanation in the string picture of QCD, with effective colour strings (tubes) connecting colour charges. In particular, the slope of both the Reggeon trajectory α'_R and the soft Pomeron trajectory α'_P are uniquely determined, being inversely proportional to the string tension κ. Thus, this parameter of the non-perturbative QCD should also enter the saturation model.

In this paper it will be demonstrated, that the string tension may is naturally built into a saturation model with the impact parameter dependence. The necessary assumption is that the impact parameter profile of the interaction is given by matter density in the target hadron. Furthermore, the parameters of the Pomeron trajectory obtained from the proposed saturation model will be shown to agree both with the results from effective string picture and with experimental values.

2. The Regge model

Let us consider elastic scattering of hadrons at c.m.s. energies squared s much larger than the typical hadronic scale s_0. We focus on the forward peak, so the momentum transfer $|t| \ll s$ $(t < 0)$. In the Regge model the elastic amplitudes are governed by singularities in the complex angular momentum plane. Assuming, that the dominant contribution to the amplitude comes from a single Regge pole, one gets

$$\mathcal{M}_R(s,t) = F_R(t) s^{\alpha_R(t)} \,, \tag{1}$$

with the linear Regge trajectory

$$\alpha_R(t) \simeq \alpha_R + \alpha'_R t \,, \tag{2}$$

and $F_R(t)$ characterising the form-factors of the colliding hadrons. At very high energies, the leading contribution to the amplitude comes from the Regge trajectory with vacuum quantum numbers — the Pomeron. The subleading Reggeons have smaller intercepts $\alpha_R < \alpha_P$, and may be neglected in the asymptotic limit $s \to \infty$.

The trajectories of subleading Reggeons may be extrapolated to positive values of t, that is from the scattering regime to the resonance regime. Families of states with masses M_i are then found along the trajectory, with spins equal to $\alpha_R(M_i^2)$. The trajectories are approximately linear, both in the scattering and in the resonance domain. Such a behaviour naturally follows from the string picture of quark interaction at large separations. The subleading Regge trajectories may be related to mesons — open string states, characterised by the value of slope

$$\alpha'_R = \frac{1}{2\pi\kappa} \,, \tag{3}$$

where the string tension $\kappa \simeq 0.18$ GeV^2.

In the effective string language, the Pomeron may be interpreted as an exchange of a closed string[1]. The resonances associated with the

[1] This statement may be motivated, for instance, by the cylindric topology of the perturbative two gluon ladder in the large N_c limit.

Pomeron should be, therefore, glueball states (see *e.g.* Ref. [14]). Recently, the Pomeron exchange was studied [15, 16] using a correspondence between QCD in four dimensions and a fundamental string theory in a special background geometry (the AdS/CFT correspondence [17, 18]). This approach gives the Pomeron slope

$$\alpha' = \frac{\alpha'_R}{4}\,, \tag{4}$$

and the Reggeon slope consistent with (3). These results reasonably agree with experimental values $\alpha'_R \simeq 0.9$ GeV^{-2} and $\alpha'_P \simeq 0.25$ GeV^{-2} [14].

Assuming the Regge form of the elastic amplitude $\mathcal{M}$, the parameters of the Pomeron trajectory can be determined,

$$\alpha_P = \frac{\partial}{\partial y} \log \mathcal{M}(s,t)\bigg|_{t=0}\,, \tag{5}$$

and

$$\alpha'_P = \frac{\partial^2}{\partial y\,\partial t} \log \mathcal{M}(s,t)\bigg|_{t=0}\,, \tag{6}$$

where we introduced a variable $y = \log(s/s_0)$ related to typical rapidity of the interaction. We consider also the slope parameter $B(s)$ defined by the relation $d\sigma^{\rm el}(s,t)/dt \simeq d\sigma^{\rm el}/dt(s,t=0)\exp(B(s)t)$. It follows that

$$B(s) = 2\,\frac{\partial}{\partial t} \log \mathcal{M}(s,t)\bigg|_{t=0}\,. \tag{7}$$

3. The saturation model

The GBW saturation model is formulated in terms of colour dipole scattering off a target proton. At high collision energies $\sqrt{s}$, the dipole size r and the impact parameter $\vec{b}$ are invariant in the interaction. The dipole–proton cross section is parametrised by $\sigma_{\rm d}(r,s) = \sigma_0[1-\exp(-r^2/4R^2(s))]$ with $R^2(s) \sim s^{-\lambda}$ and $\lambda = 0.29$. The steep rise of the cross section $\sigma_{\rm d}(r,s) \sim (s/s_0)^\lambda$ for small values of $r^2/R^2(s)$, characteristic for the hard Pomeron, saturates to a constant σ_0 at $r^2/R^2(s) \gg 1$. In this formulation, the impact parameter profile of the interaction is implicitly given by the Heaviside function, $S_0(b^2) = \theta(b_0 - b)$. We will relax this simplifying assumption, allowing more realistic profiles.

For the sake of generality, we shall neglect the hadronic wave function effects and focus on a single dipole scattering. The high energy dipole scattering amplitude is dominated by the imaginary part, thus one introduces a real function N, defined by

$$\mathcal{M}_{\rm d}(\vec{r},\vec{b},s) = isN(\vec{r},\vec{b},y)\,. \tag{8}$$

The total cross section for the dipole scattering is expressed in the standard way

$$\sigma_{\rm d}(r,s) = 2\int d^2b\; N(\vec{r},\vec{b},y)\,. \qquad (9)$$

We shall use the following form of the function N,

$$N(\vec{r},\vec{b},y) = N(\rho(r^2,y)S(b^2))\,. \qquad (10)$$

This choice is based on the assumption that scattering amplitudes at different values of the impact parameter are independent. The argument $\rho(r^2,y)S(b^2)$ may be interpreted as an amplitude for a single scattering. Thus,

$$N(\rho(r^2,y)S(b^2)) \simeq \rho(r^2,y)S(b^2) \quad \text{for} \quad \rho(r^2,y)S(b^2) \ll 1. \qquad (11)$$

On the other hand, in the blackness limit the amplitude saturates the unitarity bound

$$N(\rho(r^2,y)S(b^2)) \to 1 \quad \text{for} \quad \rho(r^2,y)S(b^2) \to \infty\,. \qquad (12)$$

Those properties are realized, for instance, by a natural extension of the GBW model incorporating a non-trivial impact parameter dependence

$$N(\vec{r},\vec{b},y) = 1 - \exp[-(r^2/r_0^2)\exp(\lambda y)S(b^2)]\,, \qquad (13)$$

The solution to the Balitsky–Kovchegov equation without impact parameter dependence suggests another choice of the unitarising function [7], $N(v) = v/(1+v)$, so one may also consider

$$N(\vec{r},\vec{b},y) = \frac{(r^2/r_0^2)\exp(\lambda y)S(b^2)}{1+(r^2/r_0^2)\exp(\lambda y)S(b^2)}\,, \qquad (14)$$

a convenient parametrisation for analytic considerations.

The features of the impact parameter profile $S(b)$ are essential in the further analysis. The presence of mass gap in QCD implies that the large values of the impact parameter should be at least exponentially supressed. The exponential tail of $S(b)$ would lead to cross section $\sigma(s) \sim \log^2 s$ saturating the Froissart bound, up to a constant factor. In fact, the Gaussian tail of the profile function is necessary to obtain the linear Regge trajectory from the saturation model. Thus, we assume

$$S(b^2) = \exp(-Ab^2)\,. \qquad (15)$$

In Sec. 4 it will be shown that this form of $S(b)$ is dictated by the properties of the confining QCD vacuum.

The elastic scattering amplitude of the colour dipole with momentum transfer $t=-q^2$ is dominated by its imaginary part,

$$\mathcal{M}(s,t=-q^2) = 2is \int d^2b\; N(\vec{r},\vec{b},y)\; \exp(i\vec{q}\vec{b}) \tag{16}$$

and consequently the elastic cross section reads

$$\frac{d\sigma_{\rm el}}{dt} = \frac{1}{4\pi}\left|\int d^2b\; N(\vec{r},\vec{b},y)\; \exp(i\vec{q}\vec{b})\right|^2 . \tag{17}$$

The slope parameter is given by

$$B(s) = \frac{1}{2}\,\frac{\int d^2b\; b^2\, N(\vec{r},\vec{b},y)}{\int d^2b\; N(\vec{r},\vec{b},y)} . \tag{18}$$

Let us introduce a useful variable

$$\xi = \rho(r^2,y) S(b^2) . \tag{19}$$

Then

$$\int d^2b\; b^2\, N(\vec{r},\vec{b},y) = \frac{\pi}{A^2} \int\limits_0^{\rho(r^2,y)} \frac{d\xi}{\xi} \log\left(\frac{\rho(r^2,y)}{\xi}\right)\; N(\xi) . \tag{20}$$

The leading term at large $\rho(r^2,y)$ can be extracted in a general case, using the fact that $N(\xi)\simeq 1$ for $\xi \gg 1$,

$$\int d^2b\; b^2\, N(\vec{r},\vec{b},y) = \frac{\pi}{2A^2} \log^2(\rho(r^2,y)) + O(\log(\rho)) . \tag{21}$$

Analogously we find

$$\int d^2b\;\; N(\vec{r},\vec{b},y) = \frac{\pi}{A} \log(\rho(r^2,y)) + O(1) . \tag{22}$$

Therefore, the slope parameter

$$B(y) = \frac{1}{4A} \log(\rho(r^2,y)) + O(1) \tag{23}$$

and

$$\alpha'_P = \frac{\partial}{\partial y} \frac{\log(\rho(r^2,y))}{8A} . \tag{24}$$

Note, that for very large rapidities, the slope is governed by the long distance behaviour of $S(b^2)$. Thus, the latter result is universal for all distributions with Gaussian asymptotics.

In the saturation model framework, $\rho(r^2, y)$ corresponds to single exchange of a perturbative pPomeron, characterised by an exponential dependence on the rapidity

$$\rho(r^2, y) = c(r^2)\exp(\lambda y)\,, \tag{25}$$

where $1 + \lambda \simeq 1.3$ is the hard Pomeron intercept and $c(r^2)$ is an increasing function of the dipole size. Therefore, we find a simple relation connecting the intercept of the *hard* (bare) Pomeron with the slope of the *soft* (unitarised) Pomeron

$$\alpha'_P = \frac{\lambda}{8A}\,. \tag{26}$$

The parameter A of the impact parameter dependence is dimensionful and one expects that it should contain information about the dimensionful parameter of QCD vacuum — the colour string tension κ. It will be suggested in the next section that, indeed, these parameters may be linked.

Assuming (25), one finds the unitarised Pomeron intercept for the Gaussian profile may be found from (22)

$$\alpha_P(y) = 1 + \frac{1}{y + \log c(r^2)}\,, \tag{27}$$

for $y + \log c(r^2) \gg 1$. The normalisation of $c(r^2)$ depends on the energy scale used in the definition of rapidity, so we specify $y = \log(s/s_0)$ with $s_0 = 1$ GeV^2. Then, the information contained in the GBW model[2] may be used to obtain $c(r^2) \simeq 1$ for large dipoles with $r \sim 1$ fm, and within the model uncertainty, $|\log c(r^2)|$ is a number of order of one. Thus, for the case of two light hadron scattering at $y \gg 1$, it is enough to retain only the leading dependence on y

$$\alpha_P(y) \simeq 1 + \frac{1}{y}\,. \tag{28}$$

The intercept slowly varies with collision energy, but for asymptotically high energies the limit $\alpha_P = 1$ is approached. It is interesting to compute the value of effective intercept for the energy range probed experimentally. Taking into account the model uncertainty, one obtains $\alpha_P = 1.11 \pm 0.01$ at $\sqrt{s} = 100$ GeV, $\alpha_P = 1.08 \pm 0.006$ at $\sqrt{s} = 500$ GeV and $\alpha_P = 1.07 \pm 0.004$ at $\sqrt{s} = 1.8$ TeV. These values are rather close to the famous value $\alpha_P = 1.08$ of the Donnachie–Landshoff Pomeron [19].

A similar analysis may be performed for the exponential profile $S(b^2) \sim \exp(-A^{(\mathrm{exp})}b)$. One finds $\alpha_P^{(\mathrm{exp})} \sim 1 + 2/y$ and ${\alpha'_P}^{(\mathrm{exp})} = \lambda^2 y/(2A^{(\mathrm{exp})})^2$. Clearly, this scenario is strongly disfavoured by the data due to the increase of α'_P with rapidity and too high values of the effective intercept.

[2] In fact, one needs to know also details of $S(b^2)$ so we anticipated the result of Sec. 4.

4. Matter distribution in the hadron

Let us consider a simple model in which a hadron is represented by a bound state of two light quarks and focus on the large distance behaviour of the wave function. In such case, only the effects of confinement are relevant and the short-range perturbative interaction may be neglected. In this approximation the Hamiltonian for two quarks, with the mass m, connected by the coloured string with the string tension κ, takes the form [20]

$$H(\vec{p}_1, \vec{p}_2, \vec{r}) = \sqrt{p_1^2 + m^2} + \sqrt{p_2^2 + m^2} + \kappa|\vec{r}| \,, \tag{29}$$

where $\vec{p}_1$ and $\vec{p}_2$ denote the quarks momenta in the meson rest frame and $\vec{r}$ is the quark–anti-quark separation vector. Note, that transverse oscillations of the string are suppressed, so that the string energy is proportional to the string length. The spectrum and eigenstates of the problem were found in Ref. [20, 21]. In the limit of very small quark masses, $m^2 \ll \kappa$ the wave function of the ground state is given by [21]

$$\psi(\vec{r}_1) = (\pi\kappa)^{-3/2} \exp(-\kappa r_1^2/2) \tag{30}$$

with the quark position with respect to the centre of mass given by the three-dimensional vector $\vec{r}_1 = \vec{r}/2$, with the longitudinal component $r_{||}$ and the transverse one $\vec{b}$. The profile of the scattering amplitude in the impact parameter plane is related to matter density in the target hadron, $S(b^2) \sim \int dr_{||} |\psi(\vec{r}_1)|^2$, thus

$$S(b^2) = \exp(-\kappa b^2) \,. \tag{31}$$

This form of the profile function explicitely depends on the string tension and will be used in the next section to reconcile the saturation model with the string picture of high energy scattering.

5. String tension and the saturation model

We arrived at the point where a quantitative comparison between the saturation model and the string model of high energy scattering can be performed. First, we recall the crucial assumptions:

A1. the unitarisation occurs at each value of the impact parameter independently;

A2. the profile of the interaction in the impact parameter is given by the matter density in the target hadron;

A3. the large distance behaviour of the hadron wave function is the same as that of two light quarks connected by a colour string.

Then, from (15), (26) and (31) one deduces that the Pomeron slope in the saturation model,

$$\alpha'_P = \frac{\lambda}{8\kappa}\,, \tag{32}$$

is a function of the string tension. This should be compared with the slope value from the string picture [15, 16]

$$\alpha'_P = \frac{1}{8\pi\kappa}\,. \tag{33}$$

Both the results depend reciprocally on the string tension and they would be identical if

$$\lambda = \frac{1}{\pi}\,, \tag{34}$$

that is $\lambda \simeq 0.32$, a value fitting well the hard Pomeron phenomenology. Conversely, using the hard Pomeron intercept from the GBW model and the string tension $\kappa = 0.18$ GeV^{-2}, one gets $\alpha'_P \simeq 0.2$ GeV^{-2}, in reasonable accordance with the experimental measurement.

Our findings may be summarised by the following form of the unitarised Pomeron trajectory at large energies:

$$\alpha_P(s,t) = \left(1 + \frac{1}{\log(s/s_0)}\right) + \frac{\lambda}{8\kappa}t\,. \tag{35}$$

This result is based on a rather general assumptions and is relevant for a wide class of models. Recall also, that the model describes the hard Pomeron exchange when the unitarity corrections are small, giving $\alpha_P^{(\mathrm{hard})} = 1+\lambda$ and ${\alpha'_P}^{(\mathrm{hard})} = 0$.

It is remarkable, that the simple saturation model reproduces well the parameters of the soft Pomeron trajectory. Let us stress, however, that in the current formulation, we strongly rely on the assumption A1 of the local saturation. Certainly, validity of this postulate may be questioned. First, the colour dipole has some size, which may contribute to the impact parameter profile of the scattering amplitude. Furthermore, some diffusion of the colour dipoles should occur and finally, the range of gluon exchange "smears out" the matter distribution in the target. A more elaborated model of colour exchange dynamics should definitely address these issues. In such a model the values of parameters may get changed and some details may get reinterpreted. Still, we expect that the crucial features of the present approach should remain. In particular, the linear Regge trajectories suggest the Gaussian form of the profile, with the width expressed by the colour string tension.

6. Conclusions

A simple saturation model incorporating an impact parameter dependence was proposed, based on the assumption of local unitarisation in the impact parameter plane. An amplitude was investigated of a colour dipole–hadron elastic scattering at small momentum transfer. The results were interpreted in terms of the Regge model, both the intercept and the slope of the Pomeron trajectory were determined. It was shown, that the Gaussian distribution of the single scattering amplitude in the impact parameter plane is preferred by the experiment.

We pointed out that the soft Pomeron may be understood either as an effective colour string exchange or as a unitarised hard (perturbative) Pomeron. The dual interpretation of the soft Pomeron implies that there exists an intimate connection between these two approaches. In particular, the colour string tension should enter the saturation model. We suggested that the colour string dynamics which governs the light hadron structure may provide the missing link. Namely, we observed that matter density in a "stringy" hadron has the Gaussian profile, whose width squared is proportional to the string tension. The exact quark-antiquark wave function was used as an input into the saturation model and consequently the reciprocal dependence of the Pomeron slope on the string tension arose. Both the intercept and the slope of the soft Pomeron trajectory were found to agree well with experimental results.

The analysis suggests that the impact parameter dependent saturation model may be applicable for soft hadronic processes. The trajectory of the soft Pomeron emerges as a result of subtle interplay between perturbative physics of the hard Pomeron and the confining properties of the QCD vacuum. It is an intriguing question whether this matching is just a coincidence or, it is rather a deep consequence of the underlying theory.

I am grateful to Professor Jan Kwieciński from whom I benefited a lot. Especially, I would like to thank him for innumerable enlightening discussions and for collaboration on many research problems, which played an important role in my understanding of particle physics.

I thank K. Golec-Biernat and R. Janik for important remarks. This paper was supported in part by the Polish State Committee for Scientific Research (KBN) grant no. 5P03B 14420.

REFERENCES

[1] L.N. Lipatov, *Sov. J. Nucl. Phys.* **23**, 338 (1976); E.A. Kuraev, L.N. Lipatov, V.S. Fadin, *Sov. Phys. JETP* **44**, 443 (1976); **45**, 199 (1977); I.I. Balitsky, L.N. Lipatov, *Sov. J. Nucl. Phys.* **28**, 822 (1978).

[2] M. Froissart, *Phys. Rev.* **123**, 1053 (1961).

[3] I. Balitsky, *Nucl. Phys.* **B463**, 99 (1996).

[4] Y.V. Kovchegov, *Phys. Rev.* **D60**, 034008 (1999); *Phys. Rev.* **D61**, 074018 (2000).

[5] E. Levin, K. Tuchin, *Nucl. Phys.* **B573**, 833 (2000).

[6] M.A. Braun, *Eur. Phys. J.* **C16**, 337 (2000).

[7] K. Golec-Biernat, L. Motyka, A.M. Staśto, *Phys. Rev.* **D65**, 074037 (2002).

[8] E. Iancu, K. Itakura, L. McLerran, *Nucl. Phys.* **A708**, 327 (2002).

[9] A. Kovner, U.A. Wiedemann, *Phys. Rev.* **D66**, 051502 (2002); *Phys. Rev.* **D66**, 034031 (2002); *Phys. Lett.* **B551**, 311 (2003).

[10] K. Golec-Biernat, M. Wüsthoff, *Phys. Rev.* **D59**, 014017 (1999); *Phys. Rev.* **D60**, 114023 (1999).

[11] N. Timneanu, J. Kwieciński, L. Motyka, *Eur. Phys. J.* **C23**, 513 (2002).

[12] J. Bartels, E. Gotsman, E. Levin, M. Lublinsky, U. Maor, *Phys. Lett.* **B556**, 114 (2003).

[13] A.I. Shoshi, F.D. Steffen, H.J. Pirner, *Nucl. Phys.* **A709**, 131 (2002).

[14] P.V. Landshoff, `hep-ph/0209364`.

[15] R.A. Janik, R. Peschanski, *Nucl. Phys.* **B586**, 163 (2000); *Nucl. Phys.* **B625**, 279 (2002).

[16] R.A. Janik, *Phys. Lett.* **B500**, 118 (2001).

[17] J. M. Maldacena, *Adv. Theor. Math. Phys.* **2**, 231 (1998).

[18] E. Witten, *Adv. Theor. Math. Phys.* **2**, 505 (1998).

[19] A. Donnachie, P.V. Landshoff, *Phys. Lett.* **B296**, 227 (1992).

[20] A.Y. Dubin, A.B. Kaidalov, Y.A. Simonov, *Phys. Lett.* **B323**, 41 (1994).

[21] E.L. Gubankova, A.Y. Dubin, *Phys. Lett.* **B334**, 180 (1994).