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MESON-DIQUARK BOSONIZATION OF QCD_2

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1. Introduction

Diquarks play an interesting role in the bosonization of the QCD action within the functional integral approach [1, 2]. Recently, it has been shown [3] that one can describe on this basis also baryons as quark-diquark bound states (for further work on path-integral bosonization see e.g. ref. [4]).

In this paper we reinvestigate diquarks in two-dimensional QCD (QCD$_2$) in the light-cone gauge using the results obtained many years ago in ref. [1]. There, meson-diquark bosonization of a coloured theory was carried out for the first time. However, in that paper the four-quark interaction term arising from gluon exchange was rearranged into $q\bar{q}$- and qq-channels somewhat artificially by introducing a so-called residual interaction.

In contrast to this we here make use of new types of Fierz identities [2] which allow one to decompose in a more natural way the interaction term into a colour singlet $q\bar{q}$- and a colour triplet qq-term avoiding any admixtures of unphysical colour octet or sixtet contributions or of a residual interaction. This and the fact that we consider unlike [2] the two-dimensional theory (in the light-cone gauge) enable us to obtain an explicit meson-diquark bosonization without any approximation.

The present paper has mainly a methodical and pedagogical aim. We want to demonstrate the powerfulness of bilocal path integral techniques in the study of quark models. The physical outcome of this new treatment of $q\bar{q}$- and qq-channels is then a modified formula for the diquark mass spectrum having other group-theoretical numerical factors as compared with the equation given in [1]. Furthermore, it becomes clear that the case of $SU(2)_c$ is no more an exceptional one, i.e. for all $N_c \geq 2$ the diquark masses are infrared divergent.

The paper is organized as follows. In sect. 2 we give a short formulation of the model and method closely following the presentation of ref. [1]. The diquark spectrum is calculated in sect. 3, and sect. 4 contains the conclusion.

2. Model and method

We start with the Lagrangian of two-dimensional QCD with exact local colour symmetry $SU(3)_c$ and, for simplicity, also with exact global flavour symmetry $SU(3)_f$

$$\mathcal{L} = -\frac{1}{2} G_{\mu\nu,\alpha\beta} G^{\mu\nu}_{\beta\alpha} + \bar{q}_{a\alpha} (i\gamma^\mu D_{\mu,\alpha\beta} - m\, \delta_{\alpha\beta}) q_{a\beta} \quad . \qquad (1)$$

Here

$$G_{\mu\nu,\alpha\beta} = \partial_\mu A_{\nu,\alpha\beta} - \partial_\nu A_{\mu,\alpha\beta} + ig\, [A_\mu, A_\nu]_{\alpha\beta}$$

is the field strength tensor,

$$A_{\mu,\alpha\beta} = \sum_{n=1}^{8} \frac{\lambda^n_{\alpha\beta}}{2} A^n_\mu$$

is the gauge field and

$$D_{\mu,\alpha\beta} = \partial_\mu \delta_{\alpha\beta} + ig\, A_{\mu,\alpha\beta}$$

is the covariant derivative; $q_{a\beta}$ is the spinor of the quark field. The indices $\alpha, \beta = 1, 2, 3$ denote colour; and $a = 1, 2, 3$, flavour. $\frac{\lambda^n}{2}$ are the generators of the colour group $SU(3)_c$ in the Gell-Mann representation.

It is convenient to consider QCD_2 in the light-cone gauge

$$A_- = A^+ = \frac{1}{\sqrt{2}} (A_0 - A_1) = 0$$

$$(a \cdot b = a_+ b^+ + a_- b^- = a_+ b_- + a_- b_+) \; .$$

In this case there exists only one independent dynamical quark variable $\hat{q} \equiv \begin{pmatrix} 0 \\ q_2 \end{pmatrix}$ and gluonic self-interactions as well as Faddeev-Popov ghosts are absent [1]. Let us consider the generating functional for Green functions of quarks

$$Z[\eta,\bar{\eta}] = C_1 \int \mathcal{D}A_+\, \mathcal{D}q\, \mathcal{D}\bar{q} \exp i\int d^2x \left\{ \mathrm{Tr}\, (\partial_- A_+)^2 \right.$$

$$\left. + \bar{q}\, (i\gamma^\mu \partial_\mu - m - g\, \gamma_- A_+)\, q + \bar{q}\eta + \bar{\eta} q \right\}$$

with

$$\gamma^{\pm} = \gamma_{\mp} = \frac{1}{\sqrt{2}} (\gamma^0 \pm \gamma^1), \quad \gamma^0 \equiv \sigma_1, \; \gamma^1 \equiv i\sigma_2, \; \gamma^5 \equiv -\sigma_3$$

and $\sigma_1, \sigma_2, \sigma_3$ being the Pauli matrices. Using the relation

$$\int d^2x \, \mathrm{Tr}\, (\partial_- A_+)^2 = -\frac{1}{2} \sum_{n,m=1}^{8} \int d^2x \, d^2y \, A_+^n(x) \, K_{nm}^{-1}(x,y) \, A_+^m(y)$$

with

$$K_{n,m}(x,y) = -i \, \delta_{mn} \, D(x-y)$$

and

$$D(x-y) = \frac{i}{\left(\partial_-^x\right)^2} \delta^{(2)}(x-y) = \int \frac{d^2k}{(2\pi)^2} D(k) \, e^{-ik(x-y)}, \qquad D(k) = \frac{-i}{k_-^2} \tag{2}$$

one obtains after integration over A_+ and rewriting the functional integral over q and $\bar{q} = q^+\gamma^0$ in terms of the variables q_2 , q_2^* with q_2^* being the Grassmann conjugate of q_2 (putting then again $q_2 \longrightarrow q$, $\eta_2 \longrightarrow \eta$ etc.)

$$\begin{aligned} Z[\eta,\eta^*] = C_2 \int \mathcal{D}q \, \mathcal{D}q^* \exp i \int d^2x \, d^2y \Big\{ & q^*(x) \, iG^{-1}(x,y) \, q(y) \\ & - \frac{i}{2} (2g)^2 \sum_{n=1}^{8} [q^*(x) \frac{\lambda^n}{2} q(x)] \, D(x-y) \, [q^*(y) \frac{\lambda^n}{2} q(y)] \\ & + [q^*(x) \, \eta(y) + \eta^*(x) \, q(y)] \, \delta^{(2)}(x-y) \Big\} \end{aligned} \tag{3}$$

with

$$\begin{aligned} G(x,y) &= \frac{-\partial_-^x}{-2\partial_+^x\partial_-^x - m^2 + i\varepsilon} \delta^{(2)}(x-y) \\ &= \int \frac{d^2k}{(2\pi)^2} \frac{ik_-}{2k_+k_- - m^2 + i\varepsilon} e^{-ik(x-y)} . \end{aligned} \tag{4}$$

To perform in (3) the integration over the quark fields one has to linearize the 4-quark interaction term. We write it in the form

$$\frac{i}{2} (2g)^2 \int d^2x \, d^2y \, q_B(y) \, q_A^*(x) \, \mathcal{K}_{AB,CD}(x,y) \, q_D(x) \, q_C^*(y)$$

with

$$\mathcal{K}_{AB,CD}(x,y) = \sum_{n=1}^{8} \frac{\lambda_{\alpha\delta}^n}{2} \frac{\lambda_{\gamma\beta}^n}{2} \delta_{il} \, \delta_{kj} \, D(x-y) \; . \tag{5}$$

Here A, B, C and D are short-hand notation for indices $A = \{i,\alpha\}$, $B = \{j,\beta\}$, $C = \{k,\gamma\}$ and $D = \{l,\delta\}$ where the first index in the brackets refers to flavour and the second one to colour. Now we rearrange the kernel $\mathcal{K}_{AB,CD}$ with the help of new types of Fierz identities proposed in [2]. For the colour group $SU(3)_c$ we use

$$\sum_{n=1}^{8} \lambda_{\alpha\delta}^n \, \lambda_{\gamma\beta}^n = \frac{4}{3} \delta_{\alpha\beta} \, \delta_{\gamma\delta} + \frac{2}{3} \sum_{\rho=1}^{3} \varepsilon_{\rho\alpha\gamma} \, \varepsilon_{\rho\beta\delta} \tag{6}$$

with $\varepsilon_{\alpha\beta\gamma}$ being the antisymmetric Levi-Civita tensor. For the flavour group $SU(3)_f$ with generators $T^a = \frac{\lambda_f^a}{2}$ one has

$$\delta_{1l}\,\delta_{kj} = \sum_{e=0}^{8} F^e_{ij}\,F^e_{kl}\,,$$

$$\{F^e,\ e = 0, 1, \ldots, 8\} = \left\{\sqrt{\tfrac{1}{3}}\,1_f,\ \sqrt{2}\,T^1,\ \ldots,\ \sqrt{2}\,T^8\right\},$$

$$\delta_{1l}\,\delta_{kj} = \sum_{g=1}^{9} H^g_{lk}\,H^g_{jl}\,,$$

$$\{H^g,\ g = 1, 2, \ldots, 9\} = \{F^e,\ e = 7, 5, 2, 0, 1, 3, 4, 6, 8\}\,.$$

Therefore, kernel (5) can be decomposed as

$$\mathcal{K}_{AB,CD}(x,y) = \left[\frac{1}{3}\,\delta_{\alpha\beta}\,\delta_{\gamma\delta}\,F^e_{ij}\,F^e_{kl} + \frac{1}{6}\sum_{\rho=1}^{3}\varepsilon_{\rho\alpha\gamma}\,\varepsilon_{\rho\beta\delta}\,H^g_{ik}\,H^g_{jl}\right] D(x-y)$$

$$= \left[\left(M^e_m\right)_{AB}\left(M^e_m\right)_{CD} - \left(M^\theta_d\right)_{AC}\left(M^\theta_d\right)_{BD}\right] D(x-y)$$

with

$$M^e_m = \sqrt{\tfrac{1}{3}}\,1_c\,F^e\,,$$

$$M^\theta_d = i\sqrt{\tfrac{1}{6}}\,\varepsilon^\rho\,H^g\,, \qquad \theta = \{\rho, g\} \tag{7}$$

(summation over repeated indices e and θ is now understood). 1_f and 1_c are the unit matrices in flavour resp. colour space.

Then, the generating functional takes the form

$$Z[\eta,\eta^*] = C_3 \int \mathcal{D}q\,\mathcal{D}q^*\,\exp i \int d^2x\,d^2y \left\{ q^*(x)\,iG^{-1}(x,y)\,q(y) \right.$$

$$+ \frac{i}{2}(2g)^2\Big[[q^*(x)\,M^e_m\,q(y)]\,D(x-y)\,[q^*(y)\,M^e_m\,q(x)]$$

$$+[q^*(x)\,M^\theta_d\,q^*(y)]\,D(x-y)\,[q(y)\,M^\theta_d\,q(x)]\Big]$$

$$\left. + [q^*(x)\,\eta(y) + \eta^*(x)\,q(y)]\,\delta^{(2)}(x-y)\right\}. \tag{8}$$

As in paper [1], the 4-quark interaction terms in (8) are then linearized by introducing bilocal fields so that the generating functional can be rewritten in the form

$$Z[\eta,\eta^*] = C_4 \int \mathcal{D}q\,\mathcal{D}q^*\,\mathcal{D}\Sigma\,\mathcal{D}\Psi\,\mathcal{D}\Psi^+$$

$$\cdot \exp i \int d^2x\,d^2y \left\{ q^*(x)\,iG^{-1}(x,y)\,q(y) + \frac{i}{2}\frac{1}{(2g)^2}\frac{\Sigma^e(x,y)\,\Sigma^e(y,x)}{D(x-y)}\right.$$

$$+ \frac{2i}{(2g)^2}\frac{\psi^{+\theta}(x,y)\,\psi^\theta(y,x)}{D(x-y)} - [q^*(x)\,M^e_m\,q(y)]\,\Sigma^e(x,y)$$

$$+ [q^*(x)\,M^\theta_d\,q^*(y)]\,\Psi^\theta(x,y) + \Psi^{+\theta}(x,y)\,[q(x)\,M^\theta_d\,q(y)]$$

$$\left. + [q^*(x)\,\eta(y) + \eta^*(x)\,q(y)]\,\delta^{(2)}(x-y)\right\}.$$

with $\mathcal{D}\Sigma \equiv \prod_e \mathcal{D}\Sigma^e$, $\mathcal{D}\Psi \equiv \prod_\theta \mathcal{D}\Psi^\theta$, $\mathcal{D}\Psi^+ \equiv \prod_\theta \mathcal{D}\Psi^{+\theta}$. The bilocal fields satisfy the relations $\Sigma^{+e}(x,y) = \Sigma^{*e}(y,x) = \Sigma^e(x,y)$ and

$\Psi^{+\theta}(x,y) = \Psi^{*\theta}(y,x) = -\Psi^{*\theta}(x,y)$ where "+" denotes the "Hermitiean" conjugate and "*" the complex conjugate.

After integration over the quark fields one finally gets

$$Z[\eta,\eta^*] = C \int \mathcal{D}\Sigma\, \mathcal{D}\Psi\, \mathcal{D}\Psi^+ \exp\{iW_{eff}[\Sigma^e,\Psi^\theta,\Psi^{+\theta}]\}\, Z[\eta,\eta^*|\Sigma^e,\Psi^\theta,\Psi^{+\theta}] \tag{9}$$

with

$$W_{eff}[\Sigma^e,\Psi^\theta,\Psi^{+\theta}] = \int d^2x\, d^2y \left\{ -i\, \mathrm{tr}\ \ln iG_\Sigma^{-1} + \frac{i}{2}\frac{1}{(2g)^2}\frac{\Sigma^e\ \Sigma^e}{D} \right.$$

$$\left. -\frac{i}{2}\, \mathrm{tr}\ \ln (1 - G_\Sigma^T\, 2M_d^\theta\, \Psi^{+\theta}\, G_\Sigma\, 2M_d^\tau\, \Psi^\tau) + \frac{2i}{(2g)^2}\frac{\psi^{+\theta}\ \psi^\theta}{D} \right\} , \tag{10}$$

$$Z[\eta,\eta^*|\Sigma^e,\Psi^\theta,\Psi^{+\theta}] = \exp\, i \int d^2x\, d^2y\, (-\eta^* G_n \eta + \frac{1}{2}\eta^* G_a \eta^* + \frac{1}{2}\eta G_a^+ \eta) . \tag{11}$$

In (11) the normal and anomalous Green functions for the quarks moving in external fields Σ, Ψ, and Ψ^+ are defined as

$$G_n = -\, iH^{-1} G_\Sigma , \qquad G_a = -\, H^{-1} G_\Sigma\, 2M_d^\theta\, \Psi^\theta\, G_\Sigma^T$$

with

$$H = 1 - G_\Sigma\, 2M_d^\theta\, \Psi^\theta\, G_\Sigma^T\, 2M_d^\tau\, \Psi^{+\tau}$$

and

$$iG_\Sigma^{-1} = iG^{-1} - M_m^e\, \Sigma^e . \tag{12}$$

G_Σ^T means the transpose of G_Σ.

In this way, QCD_2 given as a quark-gluon theory by Lagrangian (1) in the light-cone gauge has been reformulated in terms of colour singlet meson and colour triplet diquark bilocal fields. The representation (9)-(11) is explicit and exact. Notice that there is no need to introduce (unphysical) bilocal fields for colour octet or sixtet channels nor to consider residual interactions as has been done in [1]. As is shown in ref. [3], diquarks play a fundamental role in the further hadronization procedure to get baryons in addition to mesons.

3. Particle spectrum

The stationarity condition for the effective bilocal action (10) leads to the dynamical equations for the quark spectrum of QCD_2:

$$\frac{\delta W_{eff}}{\delta \Sigma^e} = 0 , \qquad \frac{\delta W_{eff}}{\delta \Psi^\theta} = 0 , \qquad \frac{\delta W_{eff}}{\delta \Psi^{+\theta}} = 0 .$$

From here one gets the following system of equations:

$$\sigma \equiv M_m^e\, \Sigma^e = -(2g)^2\, D\, (\mathrm{tr}[M_m^e\, G_n(\sigma,\Psi,\Psi^+)])\, M_m^e \tag{13}$$

$$\psi \equiv 2M_d^\theta\, \Psi^\theta = (2g)^2\, D\, (\mathrm{tr}[M_d^\theta\, G_a(\sigma,\Psi,\Psi^+)])\, M_d^\theta \tag{14}$$

$$\psi^+ \equiv 2M_d^\theta\, \Psi^{+\theta} = -(2g)^2\, D\, (\mathrm{tr}[G_a^+(\sigma,\Psi,\Psi^+)\, M_d^\theta])\, M_d^\theta . \tag{15}$$

Now we shall restrict ourselves to the consideration of the solution $\psi = \psi^+ = 0$, $G_a = G_a^+ = 0$, but $\sigma \neq 0$. In this case, using relations (12) and (4) and the fact that $\sigma_{AB} = \sigma\,\delta_{AB} = \sigma\,\delta_{ij}\,\delta_{\alpha\beta}$ as well as $G_{nAB} = G_\sigma\,\delta_{AB} = G_\sigma\,\delta_{ij}\,\delta_{\alpha\beta}$ one obtains from equation (13) in the momentum space

$$\sigma(p_-) = \frac{g^2}{\pi^2}\int dk_-\,\frac{\theta(|k_- - p_-| - \lambda)}{(k_- - p_-)^2}\int dk_+\,G_\sigma(k)\ ,$$

where

$$G_\sigma(k) = -\,G_\sigma^T(k) = \frac{ik_-}{2k_+k_- - m^2 + i\varepsilon - k_-\sigma(k)}\ ,$$

λ being an infrared cut-off parameter. This equation has been solved by 't Hooft [4]:

$$\sigma(p_-) = \frac{g^2}{\pi}\left(\frac{\mathrm{sgn}\ p_-}{\lambda} - \frac{1}{p_-}\right).$$

Next, we expand the integrand in (9) around the stationary solution σ, $\psi = \psi^+ = 0$. After the shift $\Sigma^e = \sigma\,\delta^{e0} + \Phi^e$ of the integration variable the generating functional takes the form

$$Z[\eta,\eta^*] = C\int \mathcal{D}\Phi\ \mathcal{D}\Psi\ \mathcal{D}\Psi^+\ \exp\, i\{W^\Phi_{free} + W^\Psi_{free} + W_{int}\}$$

$$\cdot\ Z[\eta,\eta^*\,|\sigma\,\delta^{e0} + \Phi^e, \Psi^\theta, \Psi^{+\theta}]$$

with

$$W^\Phi_{free} = \frac{i}{2}\int d^2x\ d^2y\ \Phi^e(x,y)\ (\Delta_\Phi^{-1}(x,y))^{ef}\ \Phi^f(y,x)\ , \qquad (16)$$

$$W^\Psi_{free} = 2i\int d^2x\ d^2y\ \Psi^{+\theta}(x,y)\ (\Delta_\Psi^{-1}(x,y))^{\theta\tau}\ \Psi^\tau(y,x)\ , \qquad (17)$$

$$(\Delta_\Phi^{-1})^{ef} = \frac{\delta^{ef}}{(2g)^2 D} - \mathrm{tr}\,(G_\sigma\, M_m^e\, G_\sigma\, M_m^f)\ ,$$

$$(\Delta_\Psi^{-1})^{\theta\tau} = \frac{\delta^{\theta\tau}}{(2g)^2 D} - \mathrm{tr}\,(G_\sigma\, M_d^\theta\, G_\sigma\, M_d^\tau)\ .$$

The explicit form of W_{int} can be found in [1]. With Δ_Φ and Δ_Ψ we denoted the propagators of the corresponding bilocal fields. These fulfil the following inhomogeneous Bethe-Salpeter equations for the scattering amplitudes in $\bar{q}q$ (colour singlet) and qq (colour octet) channels, respectively:

$$\Delta_\Phi^{ef} = (2g)^2\ D\ [\delta^{ef} + (\mathrm{tr}\,(G_\sigma\, M_m^e\, G_\sigma\, M_m^g))\ \Delta_\Phi^{gf}]\ ,$$

$$\Delta_\Psi^{\theta\tau} = (2g)^2\ D\ [\delta^{\theta\tau} + (\mathrm{tr}\,(G_\sigma\, M_d^\theta\, G_\sigma\, M_d^\rho))\ \Delta_\Psi^{\rho\tau}]\ .$$

Variation of (16) and (17) with respect to Φ and Ψ yields the homogeneous Bethe-Salpeter equations for the vertex functions of the corresponding bound states in the ladder approximation:

$$(\Delta_\Phi^{-1})^{ef}\ \Gamma_\Phi^f = 0\ , \qquad (\Delta_\Psi^{-1})^{\theta\tau}\ \Gamma_\Psi^\tau = 0\ .$$

The first equation defines the meson spectrum and has been solved by 't Hooft [5]. The second one defines the diquark spectrum. It has been investigated in [1]. In the explicit form it is given by

$$\Gamma^{\tau}_{\Psi}(p,r) = -i\ (2g)^2\ \frac{1}{3} \int \frac{d^2k}{(2\pi)^2}\ \frac{1}{k_-^2}\ G_{\sigma}(p+k)\ \Gamma^{\tau}_{\Psi}(p+k,r)\ G_{\sigma}(p+k-r)\ , \quad (18)$$

where r denotes the total momentum of the qq-pair.

For solving equation (18) it is convenient to introduce the wavefunction

$$h^{\tau}(p_-,r) = \int dp_+\ G_{\sigma}(p)\ \Gamma^{\tau}_{\Psi}(p,r)\ G_{\sigma}(p-r)$$

and to rewrite (18) as an equation of 't Hooft's type [5,1]

$$\nu^2\ h^{\tau}(x) = \left(\frac{\beta}{x} + \frac{\beta}{1-x} \right) h^{\tau}(x) - P\int_0^1 \frac{h^{\tau}(y)}{(y-x)^2}\, dy\ . \quad (19)$$

Here $x = p_-/r_-$, P denotes the principal value and

$$\beta = 3\pi\, \frac{m^2}{g^2} - 3\ ,$$

$$\nu^2 = \frac{3\pi}{g^2} \left[2r_+r_- - \frac{4g^2}{3\pi\lambda}\, |r_-| \right] .$$

The solution of (19) has to fulfil the boundary conditions $h^{\tau}(x) \sim x^{\gamma}$ $[(1-x)^{\gamma}]$ for $x=0$ $[x=1]$ where $\pi\gamma\cos\pi\gamma = -\beta$. The

system of eigenfunctions $h^{\tau}_k(x)$ is complete and orthogonal. For large k one obtains [5]

$$h^{\tau}(x) \cong \sqrt{2}\,\sin\pi kx\ , \quad k \gg 1\ , \quad \nu_k^2 \simeq \pi^2 k\ .$$

The diquark mass spectrum is then given by

$$M_k^2 = (2r_+r_-)_k = \frac{g^2}{3\pi}\,\nu_k^2 + \frac{4g^2}{3\pi\lambda}\,|r_-|\ . \quad (20)$$

In this way, the diquark masses go to infinity for $\lambda \longrightarrow 0$ analogously to the quark masses . This result coincides with the fact that coloured states like quarks or diquarks are not observable - they are confined. The group-theoretical numerical factors in (20) differ from those of ref. [1].

4. Conclusion

The main result of this paper is an exact meson-diquark bosonization of QCD_2 obtained by applying bilocal path integral techniques. In the present essentially improved approach there was no need to consider any residual interaction. And no unphysical colour octet or sixtet fields like in [1] have been introduced for which no attraction, and therefore, no bound states exist.

The use of new types of Fierz identities leads to a mass formula for diquarks which in the group-theoretical numerical factors differs from the old result in [1]. Although we considered here explicitly only the three colour case, it becomes clear that one now obtains an infrared divergent diquark mass formula also for $SU(2)_c$ in contrast to the earlier work.

In a forthcoming paper, the Bethe-Salpeter equation for baryons will be investigated within this model.

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Эберт Д., Кашлун Л. E2-90-374
Мезон-дикварковая бозонизация для $КХД_2$

Вновь исследуются дикварки в двухмерной КХД для случая трех цветов. В калибровке светового конуса $КХД_2$ бозонизируется с помощью билокальных функциональных методов в явной форме и без каких-либо приближений. При этом, возникающее из глюонного обмена, выражение для 4-кваркового взаимодействия преобразуется с помощью такого тождества Фирца, которое включает только синглетные по цвету кварк-антикварковые и триплетные по цвету кварк-кварковые каналы.

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Ebert D., Kaschluhn L. E2-90-374
Meson-Diquark Bosonization of QCD_2

Diquarks in two-dimensional QCD for the three colour case are reinvestigated. By means of bilocal functional techniques QCD_2 in the light-cone gauge is bosonized explicitly and without any approximation. Thereby, the four-quark interaction term arising from gluon exchange is Fierz rearranged into a form which involves only colour singlet quark-antiquark and colour triplet quark-quark channels.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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