A simple lattice model inducing a gauge theory is considered. The model describes an interaction of a gauge field to an $N \times N$ complex matrix scalar field transforming as a field in the fundamental representation. In contrast to the Kazakov-Migdal model the model contains only the linear interaction between scalar and gauge lattice fields. This model does not suffer from extra local $U(1)$ symmetries.

In an approximation of a translation invariant master field the large $N$ limit of the model is investigated. At large $N$ the gauge fields can be integrated out yielding an effective theory describing an interaction of eigenvalues of the master field. The reduced model exhibits phase transitions at the points $\beta_{\text{cr}}$ and $\beta_{\bar{\text{cr}}}$ and the region $(\beta_{\text{cr}}, \beta_{\bar{\text{cr}}})$ separates the strong and weak regions of the model.

To study the behaviour of the model at large $N$ in more systematic way the quenched momentum prescription with constraints for treating the large $N$ limit of gauge theories is used. With the help of the technique of orthogonal polynomials nonlinear equations describing the large $N$ limit of the reduced model with quenching are presented.

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1 Introduction

This paper is inspired by recent studies of the Kazakov Migdal (KM) model \[1\] - \[14\], which was suggested as a possible model of hadrons. The KM model contains a scalar field Φ in the adjoint representation of the $SU(N)$ group interacting with the usual lattice gauge field. The main point which allows to study the large N limit of the model analytically, is the absence of the kinetic term for the gauge field.

There are several problems with the KM model. A relation of this model with QCD and, in particular, the property of asymptotic freedom still should be clarified. The KM model has an extra local gauge symmetry \[9, 11\], which should be broken if we want to recover QCD from this lattice model. There is also a more technical problem, concerning the method of investigation of the large N limit in the model. Most of previous considerations used the approximation of the constant master field. To clarify this assumption the large N limit of the KM model has been investigated by an other method, namely, by the quenched momentum prescription \[15\] - \[23\] for study the large N limit of gauge theories \[24\]. It has been shown that using the quenched procedure with constraints one obtains the one side model with the quartic interaction, that makes the model not exactly solvable.

In this paper we are going to construct a modified model which does not suffer from extra gauge symmetry and admits more simple analytical treating at large $N$.

Several modifications of the KM model have been proposed recently. Migdal \[12\], and Khokhlachev and Makeenko \[13\] proposed to induce QCD by fermions in the adjoint representation of the gauge group $SU(N)$ with various types of lattice fermions. They discussed also adjoint scalar and fermionic models at large number of flavours $N_f$.

We will consider a simple lattice gauge model in which an interaction of gauge fields is induced by a scalar field, being an $N \times N$ complex matrix transforming as a field in the fundamental representation. In contrast to the KM model the modified model contains only linear interaction between scalar and gauge lattice fields. This model does not suffer from extra local $U(1)$ symmetries. To estimate the behaviour of this model we perform an integration over lattice gauge fields and then use a translation invariant master field to calculate the remaining integral over scalar field. The integration over gauge fields at large N is performed by using the Brezin-Gross results \[26\] of study one link gauge model in an external matrix field. The phase transition which takes place at large N for one link gauge model in an external matrix field \[27, 26\] gives rise to a non-trivial phase structure of our model. The reduced model exhibits phase transitions at the points $\beta_{cr}$ and $\beta_{\overline{cr}}$ and the region $(\beta_{cr}, \beta_{\overline{cr}})$ separates the strong and weak coupling regimes of the model.

To study the behaviour of the model at large $N$ in a more systematic way the quenched momentum prescription with constraints for treating the large N limit of gauge theories is used. This prescription leads to a quadratic dependence of the reduced action on unitary matrix $V$ instead of a quartic dependence for the KM model. This quadratic dependence permits to use the technique of orthogonal polynomials \[29, 1\] and to obtain nonlinear equations describing the large N limit of the reduced model with quenching. These equations are similar to equations describing the KM reduced model without quenching.
2 The Action of the Model

Let us consider a simple lattice model which describes an interaction between $N \times N$ matrix scalar field $\Phi$ and gauge field and $U_\mu$. The partition function is given by

$$Z = \int \prod_{x,\mu} dU_\mu(x) \prod_x d\Phi(x) \exp\{S(\Phi(x), U_\mu(x))\}.$$ \hspace{1cm} (1)

Where $\Phi$ is an arbitrary $N \times N$ matrix and the scalar field $\Phi(x)$ defined on the sides of a d-dimensional rectangular lattice and $U_\mu$ is an unitary matrix defined on the links; $d\Phi = \prod_{i,j} d\phi_{ij}$ and $dU_\mu$ is the Haar measure on the group of unitary matrices. The action has the following form

$$S(\Phi, U) = -\frac{1}{2} m_0^2 N \text{tr} \Phi \Phi^\dagger + \beta N \sum_{\mu > 0} \text{tr} [\Phi(x) U_\mu(x) \Phi^\dagger(x + \mu) + \Phi(x + \mu) U_\mu^\dagger(x) \Phi^\dagger(x)].$$ \hspace{1cm} (2)

As in the case of the KM model in (2) there is no a kinetic term for the gauge field $U_\mu$.

The action (2) is invariant under gauge transformations

$$U_\mu(x) \to \Omega(x) U_\mu(x) \Omega^\dagger(x + \mu), \quad \Phi(x) \to \Omega(x) \Phi(x) \Omega^\dagger(x), \quad \Phi^\dagger(x) \to \Phi^\dagger(x).$$ \hspace{1cm} (3)

Note that one can consider the matrix field $\phi_{ij}$ as a field with upper and down indices $\phi_{ij} = \phi^j_i$ and regard the upper index as the usual colour index and the down as the flavour index. So, with the respect to the down indeces the scalar field transforms as the usual matter field in the fundamental representation, and the action is the sum of $N$ copies of matter actions in the fundamental representation. In accordance with this analogy the action is invariant under the global $U(N)$ symmetry

$$\Phi(x) \to G^\dagger \Phi(x), \quad \Phi^\dagger(x) \to \Phi^\dagger(x) G;$$ \hspace{1cm} (4)

$G$ belongs to $U(N)$. Note also that in contrast to the KM model the model (2) does not have extra local $U(1)$ symmetries.

Calculating gaussian integral over fields $\Phi$ and $\Phi^\dagger$ one rewrites the model (2) as a pure gauge theory with the following reduced action

$$S_{\text{ind}} = N \sum_{\Gamma} \left( \frac{m_0^2}{\beta} \right)^{-(\ell(\Gamma))} \text{tr} U(\Gamma),$$ \hspace{1cm} (5)

Notice that the expansion (5) starts with the standard Wilson action. Recall that the analogous expansion for the KM model starts with the square of the Wilson action. To get the Wilson action instead of the action (5) one has to find a mechanism of suppression of long loops in the effective action (3). For example, one can use the following trick, which was also used by Khokhlov and Makeenko [13] to the KM model. Let attache an extra index to the field $\Phi$, i.e. introduce the field $\Phi^f$, $f = 1, ..., n_f$ and consider $m_0 \to \infty$ and $n_f \to \infty$ so that $m_0^2 \sim (n_f)^{1/4}$. Then only the single-plaquette survives in (5), the longer loops are suppressed at least as $n_f^{-1/2}$.

A naive continuous version of our model corresponds to $N$ copies of the gauged scalar fields in the fundamental representation. These $N$ copies can be interpret as fields with
different flavour and $N_c = N_f$. The corresponding Lagrangian has no an explicit Yang-Mills term

$$\mathcal{L} = \frac{N}{g_0} (\partial \Phi - i A_\mu \Phi^\dagger) (\partial \Phi^\dagger) + i A_\mu \Phi^\dagger + V(\Phi \Phi^\dagger). \quad (6)$$

Here there is not the problem which arises in the naive continuous version of the KM model, where a commutator $[A_\mu, \Phi]$ presents and therefore we miss the diagonal degrees of freedom of the gauge field.

### 3 Reduced Model

In this section we estimate the partition function (1) in the large $N$ limit using the quenched momentum prescription [15] - [23]. This prescription for treating the limit of infinite $N$ of matrix field theories has been obtained more than ten years ago by Eguchi and Kawai [15], Bhanot, Heller and Neuberger [16], Parisi [17], and Gross and Kitazawa [18] (see also [19] - [23]). Generally speaking, according to this prescription one gets the vacuum energy in the large $N$ limit by integration the free energy of the quenched reduced model over the quenching parameter $p$. The reduced model is described by a reduced action and to get the reduced action one should replace a matrix field $\Phi(x)$ with

$$D(x) \Phi^\dagger(x),$$

where $D(x) = e^{ip_\mu x_\mu}$, and $p_\mu$ is the diagonal matrix with matrix elements $p_{i\mu}, i = 1, \ldots N$. (quenching parameters). It has been shown that quenched theory produces the standard Feynman diagrams for invariant Green functions in all orders in perturbation theory.

Performing the reduction procedure $\Phi(x + \mu) = D_\mu \Phi(x) D_\mu^\dagger$ and a change of variables $U_\mu \to U'_\mu = U_\mu D_\mu$ we left with an action

$$S(\Phi, U, D) = -\frac{1}{2} m^2 N \text{tr} \Phi \Phi^\dagger + \beta N \sum_{\mu > 0} \text{tr} [\Phi U_\mu \Phi^\dagger D_\mu + \Phi U_\mu^\dagger \Phi^\dagger D_\mu] \quad (7)$$

Integrating over $\Phi$ and $\Phi^\dagger$ one gets the sum over $\Gamma$ of terms $m^{-2l(\Gamma)} \text{tr} U(\Gamma) \text{tr} D(\Gamma))$. Since all $D$ are diagonal matrices for the closed contour $\Gamma$ we have $D(\Gamma) = I$, and, therefore, we can neglect the last multipliers in both terms in equation (7) and we left with the following reduced action

$$S_{red}(\Phi, U) = -\frac{1}{2} m^2 N \text{tr} \Phi \Phi^\dagger + \beta N \sum_{\mu > 0} \text{tr} [\Phi U_\mu(x) \Phi^\dagger + \Phi U_\mu^\dagger \Phi^\dagger] \quad (8)$$

An other way of obtaining the reduced action (8) consists in application of the reduced procedure directly in the induced action (5). Performing the reduction procedure for the gauge field $U_\mu(x + \nu) = D_\nu U_\mu(x) D_\nu^\dagger$ we get the sum of terms $U(\Gamma)$, this answer follows also from (8). The same action gives the approximation of the translation invariant master field.

Hence, we have to estimate at large $N$ the following free energy

$$F = \ln \int d\Phi dU_\mu \exp \left((-\frac{\bar{m}^2}{2} N \text{tr} \Phi \Phi^\dagger + \bar{\beta} N \sum_{\mu > 0} \text{tr} \Phi (U_\mu + U_\mu^\dagger) \Phi^\dagger)\right) \quad (9)$$

where $\bar{\beta} = a^d \beta$, $\bar{m}^2 = a^d m^2$. 
Note that in (9) \( \Phi \) is an arbitrary complex \( N \times N \) matrix. Any matrix \( \Phi \) may be written in the form \( \Phi = PV \) with \( P \) hermitian and \( V \) unitary. If we perform the change of variables \( U \to S^\dagger V U^\dagger S \) in the integral

\[
I(\Phi, \Phi^\dagger) = \int dU_\mu \exp(\beta N \sum_{\mu>0} \text{tr} \Phi (U_\mu + U_\mu^\dagger) \Phi)
\]

it is then clear that the integral (10) depends only upon \( \Lambda \). Up to a normalization factor the measure \( d\Phi \) has the following form

\[
d2N^2\Phi = \Delta^2(\lambda) \prod_i d\lambda_i ds dV.
\]

The integral (10) defines the well known object, this is the partition function of the Brezin-Gross model which describes one link gauge field in the external matrix source \([25, 26]\). The form of this partition function in large \( N \) limit depends on the magnitude of external source. In our case this parameter is

\[
S = \frac{1}{\sqrt{\beta}} \sum_i |\phi_i|^2 - 2
\]

where \( \phi_i = (\Lambda)_{ii} \). The weak coupling regime corresponds to \( S < 2 \), and the strong coupling regime to \( S > 2 \). The answer for both regimes can be written in the following form \([26]\)

\[
\frac{F}{N^2} = \frac{2}{N} \sum (\beta^2 |\phi|^4 + c)^{1/2} - \frac{1}{2N^2} \sum_{i,j} \log[(\beta^2 |\phi|^4 + c)^{1/2} + (\beta^2 |\phi|^4 + c)^{1/2}] - (c + 3/4)
\]

The constant \( c \) is different for different regimes. For the weak coupling regime \( c = 0 \) \([25]\) and for the strong coupling regime the constant is determined from the equation

\[
\frac{1}{N} \sum (\beta^2 |\phi|^4 + c)^{-1/2} = 2
\]

The next step is the calculation the integral in respect of \( \phi \)'s

\[
F = \ln \int \prod_i (\phi_i - \phi_j)^2 \prod d\phi_i d\phi_j \exp[\theta(S - 2)\delta(c) + \theta(2 - S)\delta(c - 1)]
\]

\[
\exp[-\frac{m^2}{2} N \sum \phi_i^2 + 2Nd \sum (\beta^2 |\phi|^4 + c)^{1/2} - \frac{d}{2} \sum \log[(\beta^2 |\phi|^4 + c)^{1/2} + (\beta^2 |\phi|^4 + c)^{1/2}] - (c + 3/4)]
\]

### 3.1 Week Coupling Reduced Model

One can roughly estimate the behaviour of this model in week coupling regime ignoring the contribution from the second term in the square brackets in the first line of (14) and substituting in (14) an average of

\[
<\phi_j^2> = \frac{1}{N} \text{tr} \Phi^2 = \frac{1}{\mu}
\]

instead of the sum over \( j \). In this case we get the usual 1-matrix problem in the logarithmic potential

\[
\mathcal{V} = -\frac{m^2}{2} N \text{tr} \Phi^2 + \frac{d}{2} \text{tr} \log(1 + \mu \Phi^2)
\]

(16)
where \( m_0^2 = 4d \bar{\beta} - \bar{m}^2 \). The parameter \( \mu \) should be defined from the selfconsistent equation (15). Taking into account only quadratic terms in (16) we get the following selfconsistent equation

\[
d\mu - m_0^2 = \mu,
\]

(17)

which has a solution for \( d > 1 \) only if \( m_0^2 > 0 \), i.e. \( \bar{\beta}_{cr} = \bar{m}^2/4d \) is a critical point (the symbol \( \bar{cr} \) means that this critical point is obtained starting from week coupling regime). There is also a restriction on \( d < 5/4 \) coming from the requirement of applicability of the used approximation. According to this requirement the region \( |\phi| < 2/\sqrt{\mu} \) must be inside the region where the potential (16) increases, i.e. \( 4/\mu < d/m_0^2 - 1/\mu \).

However there is a difficulty which makes the application of the quadratic potential \( V(\Phi) \) in the weak coupling regime non-consistent. It is caused by a breaking of the equality \( S < 2 \). This gives a raison to add a quartic term to the potential

\[
V(\Phi) = \frac{\bar{m}^2}{2} N \text{tr} \Phi^2 + \bar{g} N \text{tr} \Phi^4
\]

(18)

Here \( \bar{m}^2 < 4d \bar{\beta} \), and, therefore, \( m_0^2 \) is positive, i.e. we have the Goldstone potential; \( \bar{g} = a^d g \). In this case equation (23) gives \( \mu = \frac{2m_0^2}{2d+1} \). In the adopted approximation the condition \( S < 2 \) means \( \frac{\bar{m}^2 g_{eff}}{g_{eff}} > 2 \), that can be achieved for \( \beta \) being large enough; \( g_{eff} = \bar{g} - \frac{d}{4} \mu \). This approximation has also the following restriction \( \frac{\bar{m}^2 g_{eff}}{g_{eff}} > \frac{8(1+2d)}{m_0^2} \). Note that all the numerical factors are obtained in the approximation of the quadratic potential in the vicinity of the minimum of the Goldstone potential.

One should note that the application of the reduction procedure itself for the weak coupling regime is not quite correct since in this region as it was mentioned in the Introduction one has to use the reduction with constrains (see [24]). However the simple calculations presented above show that we have to deal with the effective Goldstone potential and this gives an indication that the structure of the model with quenching momentum may be similar to the phase structure of spin glasses [19].

### 3.2 Strong Coupling Reduced Model

Now we are going to estimate the free energy (21) in the strong coupling regime. In this case the second term in the first line of (24) gives a main contribution, i.e. we have to deal with (22), where \( c \) satisfies (23). Taking the following approximation for (22)

\[
F = -Ng_{eff} \beta^2 \sum \phi^4 + \text{const}, \quad g_{eff} = \frac{1 + 2\sqrt{c}}{2\sqrt{c}(1 + \sqrt{c})}
\]

(19)

we get the one matrix problem in quartic potential with a wrong sign. A phase transition in this model occurs at \( \bar{\beta}_{cr} = \frac{1}{3} \sqrt{\frac{5}{3}} \sqrt{\frac{m_0^2}{d}} \), the point where the function of distribution of eigenvalues lose the positivity (here \( \bar{cr} \) means that this critical point is obtained starting from strong coupling regime). It is interesting to compare the strong coupling regime of (9) with the strong coupling regime of the KM model. We will do it in the forthcoming publication.

The above estimations advocate for the following picture. There are two phase in the theory (1), weak and strong phases. They are separated by the intermediate region \( (\bar{\beta}_{cr}, \bar{\beta}_{cr}) \). The presented estimations are rather rough and in principle they can be made
more precise if one does not restrict by the quartic potential and deal with logarithmic potential. However we believe that the qualitative picture of the phase structure of the model (9) is correctly reproduced. The condition $\beta_{cr} < \beta_{cr}$ gives a restriction on $d$.

4 Reduced Model with Quenching

In the case of gauge theories the quenching of the momentum must be accompanied by a constraint on the eigenvalues of the covariant derivative. Without this constraint one gets naive reduced model without quenching. In the case of the Wilson gauge theory the naive reduce model describes correctly the theory only in the strong coupling regime. The suitable constraint restricts the eigenvalues of $U_\mu D_\mu$ to be equal to $D_\mu$. This means that we have to take the measure $d\mu(U)$ in (9) to be

$$d\mu(U) = \prod \mu C(U_\mu, D_\mu)$$

where $dU_\mu$ is the Haar measure on $SU(N)$ and

$$C(U_\mu, D_\mu) = \prod \mu \int dV_\mu \delta(U_\mu - V_\mu D_\mu V_\mu^\dagger D_\mu^\dagger)$$

If we integrate out the $U_\mu$, setting $U_\mu = V_\mu D_\mu V_\mu^\dagger D_\mu^\dagger$ we obtain

$$\exp(\tilde{F}(p)) = \int d\Phi \prod \mu dV_\mu \exp\{d\left(-\frac{1}{2}m^2 N \text{tr} \Phi \Phi^\dagger + \beta N \sum_{\mu>0} \text{tr} \left[\Phi V_\mu D_\mu V_\mu^\dagger D_\mu^\dagger + \Phi V_\mu D_\mu V_\mu^\dagger \Phi^\dagger D_\mu^\dagger \right]\right)\}$$

To find the vacuum energy at the large $N$ limit one should integrate $\tilde{F}(p)$ over $dp_\mu$. In fact we will use the formula

$$F = \int d\mu(p) \tilde{F}(p),$$

where

$$d\mu(p) = \prod_i \frac{d^d p_i}{(\sqrt{\pi}a)^d} \exp(-p_i^2 a^2).$$

So, the reduced action with quenching for the model (2) contains the quadratic dependence from the unitary matrix $U$. Note that the reduced action with quenching for the KM model has the quartic dependence from the unitary matrix, that makes the model not exactly solvable.

The first step in the calculation of (22) is the evaluation of the integral over $V_\mu$’s. It is interesting to note that just for this case the technics developed by Kazakov and Migdal [1, 2], and Gross [11] permits to study analytically the problem in question.

Using the two-matrix representation of the Itzykson-Zuber (IZ) integral [28] and the orthogonal polynomial method [29] the following integral over $V_\mu$’s

$$\exp I(\{\phi^2\}, \{p_i\}) = \int dV \exp(N \text{tr} \Phi (V(D + D^\dagger)V^\dagger \Phi^\dagger))$$

at large $N$ can be represented as

$$\exp I(\{\phi^2\}, \{p_i\}) = N^2 \int_0^1 dt (1 - t) \ln f(t),$$
where
\[ f(t) = -t + \oint \frac{dz}{2\pi i} \frac{V'_1(p(z), t)}{z^2}, \]
\[ p(y, t) = \frac{1}{y} + \oint \frac{dz}{2\pi i} \frac{V'_2(p(z), t)}{1 - zy f(t)} \]
\[ q(y, t) = \frac{1}{y} + \oint \frac{dz}{2\pi i} \frac{V'_1(p(z), t)}{1 - zy f(t)} \]
and
\[ V_1(x) = \frac{1}{N} \sum_b \ln(x - \phi_b^2), V_2(y) = \frac{1}{N} \sum_b \ln(y - 2 \cos ap_b), \] (28)

We left only linear dependence on \( D \) and \( D^\dagger \) in the action in (25), since the dropped \( D' \)s do not contribute to closed loops. Note that in spite of this reason such step needs more arguments.

The next step is the evaluation of the integral over \( \Phi \) and \( p_\mu \). Using the replica trick \[17, 19\]
\[ \ln \int d\phi \Delta^2(\phi) \exp\{I(\{\phi^2\}, \{p_i\})\} = \lim_{n \to 0} \frac{1}{n} \left( \int dp \prod_{a=1}^n d\phi^a \Delta^2(\phi^a) \exp\{I(\{\phi^a\}^2), \{p_i\})\} - 1 \right) \] (29)

one can perform first the integration over \( p_\mu \)’s and then use the saddle-point approximation for integral over \( \phi \)
\[ 2 \sum_{j \neq i} \frac{1}{\phi_i^2 - \phi_j^2} = 2N \frac{m^2_{\phi^2}}{\beta} \phi_i^2 - d \frac{\partial}{\partial \phi_i} J(\{\phi^a\}) \] (30)

where
\[ \ln J(\phi) = \int dp \exp I(\{\phi^a\}^2), \{p_i\}), \] (31)

\( \alpha \) is the replica index. Equation (30) is more complicated then the corresponding equation for the reduced KM model and at present we do not know how to solve it. However, an obvious advantage of the model (1) is that the IZ integral completely describes the behaviour of the model, while this integral describes the KM model only in the approximation without quenching, and even exhaustive knowledge of the result of integration of the IZ integral over eigenvalues of \( \Phi \) does not provide information about the behaviour of the KM model in all regions of coupling constant.

### 5 Concluding Remarks

A lattice gauge model with simple linear interaction between gauge and scalar fields was considered. This model does not suffer from local \( U(1) \) symmetries and it contains the Wilson action as the first term in the expansion of the effective action. So, it seems the model can be considered as a rather suitable approximation to QCD. It was shown that the model in the large \( N \) limit admits the analytical investigation and has the rich phase structure. There are many open problems with the model. The most important question here, as well as for the KM model, is a clarification of the status of asymptotic freedom.

There are also some more technical problems which must be tried out. First, it would be interesting to study the reduced model in an external random field. The action
of this model is given by the expansion (7). Note that the role of the random field \( D \) consists in the suppression of contributions coming from open loops. A similar suppression may be ensured by abelian random fields \( \exp(i\varphi_\mu) \). Such model will be a subject of our forthcoming paper [31]. Another technical questions, as it has been mentioned in Section 3, are related with more thorough estimation of the free action (14).

In concluding let us note once again that though there are many open questions with induced QCD, this subject seems to be very promising and undoubtedly is worthy of further developments.

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