acta physica slovaca vol. 49 No. 2, 231 – 239 **April 1999**

PHONON-LIKE EXCITATIONS OF THE INSTANTON LIQUID*

S.V. Molodtsov†[1], A.M. Snigirev‡[2], G.M. Zinovjev¶$[3]

† *State Research Center, Institute of Theoretical and Experimental Physics, 117 259 Moscow, Russia*

‡ *Nuclear Physics Institute, Moscow State University, 119 899 Moscow, Russia*

¶ *Fakultät für Physik, Universität Bielefeld, 335 01 Bielefeld, Germany*

$ *Bogolyubov Institute for Theoretical Physics, National Academy of Sciences of Ukraine 252 143 Kiev, Ukraine*

Received 2 October 1998, accepted 1 December 1998

The phonon-like excitations of (anti)-instanton liquid ($\bar{I}I$) due to adiabatic variations of vacuum wave functions are studied in this paper. The kinetic energy term is found and proper effective Lagrangian for such excitations is evaluated. The properties of their spectrum, while the corresponding masses are defined by Λ_{QCD} with prevailing chromoelectric component, are investigated based on the phenomenology of QCD vacuum already developed.

PACS: 11.15.Kc, 12.38.-t, 12.38.-Aw

Today's great splash of interest in instanton-based models was provoked by recent attempts to advance our understanding of QCD behaviour at very high densities [1]. It was argued there that instantons are able to generate a gap quite suitable to revive the mechanism of the 'colour superconductor' in QCD for quantitative utility. In fact, the instantons and, in particular, the model of (anti)-instanton liquid correctly seize many nonperturbative phenomena and vacuum features as chiral symmetry breaking, gluon condensate, topological susceptibility [2],[3]. However, pushing forward the detailed comprehension of QCD principal issues, new questions appear that can be addressed to the instanton liquid model for its own sake. We are focused on one of those in this paper.

It is usually supposed the corresponding functional integral in this approach is saturated by quasi-classical configurations close to the exact solutions of the Yang-Mills equations (the Euclidean solutions called the (anti)-instantons) and the wave function of vacuum being homogeneous in metric space is properly reproduced by averaging over their collective coordinates.

*Dedicated to Ján Pišút on the occasion of his 60th birthday.

[1] E-mail address: molodtsov@vitep5.itep.ru

[2] E-mail address: snigirev@lav.npi.msu.su

[3] E-mail address: gezin@hrz.uni-bielefeld.de, gezin@ap3.bitp.kiev.ua

In the $\bar{I}I$ liquid approach one takes the superposition ansatz of the pseudo-particle (PP) fields as a relevant approximation to the 'genuine' vacuum configuration

$$A_\mu(x) = \sum_{i=1}^{N} A_\mu(x;\gamma_i) \ . \tag{1}$$

Here $A_\mu(x;\gamma_i)$ denotes the field of a singled (anti)-instanton in singular gauge with $4N_c$ (for $SU(N_c)$ group) coordinates $\gamma = (\rho,\ z,\ \Omega)$, of the size ρ with the coordinate of its centre z and Ω as its colour orientation. Then,

$$A^a_\mu(x;\gamma) = \frac{2}{g}\ \Omega^{ab}\bar{\eta}_{b\mu\nu}\frac{y_\nu}{y^2}\frac{\rho^2}{y^2+\rho^2}\ ,\ \ y = x - z\ , \tag{2}$$

where η is the 't Hooft symbol [4]. For anti-instanton $\bar{\eta} \to \eta$ (making the choice of the singular gauge allows us to sum up the solution preserving the asymptotic behaviour). For simplicity we shall not introduce the different symbols for instanton and anti-instanton, and then in the superposition of Eq. (1) N becomes the total number of PP in the system 4-volume V with the density $n = N/V$.

The 'liquid' is introduced by the distribution function over instanton sizes and in one-loop approximation it is taken as the following

$$\mu(\rho) = \rho^{-5}\widetilde{\beta}^{2N_c}e^{-\beta(\rho)-\nu\rho^2/\overline{\rho^2}}\ . \tag{3}$$

As known, the respective parameters are defined by the QCD renormalization scale $\Lambda_{\rm QCD}$ together with the constant ξ characterizing effective infrared repulsion of the PP's in the model of Ref. [3], for example,

$$\left(\overline{\rho^2}\right)^2 = \frac{\nu}{\beta\xi^2 n}\ ,\ \ \xi^2 = \frac{27}{4}\frac{N_c}{N_c^2-1}\pi^2\ ,\ \ \nu = \frac{1}{2}(b-4)\ ,\ \ b = \frac{11}{3}N_c\ , \tag{4}$$

with the Gell-Mann–Low beta function

$$\beta(\rho) = -\ln C_{N_c} - b\ln(\Lambda\rho)\ ,\ \ \Lambda = \Lambda_{\overline{MS}} = 0.92\Lambda_{P.V.}\ , \tag{5}$$

and constant C_{N_c} depending on the regularization scheme, in particular,

$$C_{N_c} \approx \frac{4.6\ \exp(-1.37N_c)}{\pi^2(N_c-1)!(N_c-2)!}\ .$$

The parameters β and $\widetilde{\beta}$ are fixed at the characteristic scale $\bar{\rho}$, i.e.

$$\beta = \beta(\bar{\rho})\ ,\ \ \widetilde{\beta} = \beta + \ln C_{N_c}\ .$$

The distribution $\mu(\rho)$ has obvious physical meaning, namely, the quantity $d^4x\ d\rho\ \mu(\rho)$ is proportional to the probability to find an instanton of size ρ in some point of a volume element d^4x. At small ρ the behaviour of distribution function is stipulated by the quantum-mechanical uncertainty principle preventing a solution being compressed at a point (radiative correction). At large ρ the constraint comes from repulsive interaction between the PP's which is amplified with (anti)-instanton size increasing.

In what follows, we study the excitations of $\bar{I}I$ liquid induced by the adiabatic dilatational deformations of instanton solutions. We show the corresponding excitations display a particle-like ('phonon') behaviour, developing the mass spectrum settled by $\Lambda_{\rm QCD}$ and with the chromoelectric component prevailing. It looks quite plausible that incorporating light quark condensate would allow us to describe the light hadrons on the correct mass scale, and to provide more insight into the confinement mechanism.

The guiding idea of selecting deformation originates from transparent observation. The deformations measured in the units of action $\frac{dq\ dp}{2\pi\ \hbar}$ (here $q,\ p$ are the generalized coordinate and momentum) have only physical meaning. However, the instantons are characterized by 'static' coordinates γ and, therefore, need to appoint the momenta conjugated. It looks quite natural for the variable ρ to introduce $\dot{\rho} = d\rho/dx_4$ as a conjugated momentum. Then we should include additional degrees of freedom induced by the variation of instanton size, and under sufficiently slow variation of its size the additional terms occur in strength gluon tensor

$$G'^a_{\mu\nu} = G^a_{\mu\nu} + g^a_{\mu\nu}\ . \tag{6}$$

The first term corresponds to the contribution generated by the instanton profile

$$G^a_{\mu\nu} = -\frac{8}{g}\frac{\rho^2}{(y^2+\rho^2)^2}\left(\frac{1}{2}\bar{\eta}_{a\mu\nu} + \bar{\eta}_{a\nu\rho}\frac{y_\mu y_\rho}{y^2} - \bar{\eta}_{a\mu\rho}\frac{y_\nu y_\rho}{y^2}\right)\ ,$$

and in adiabatic approximation the corrections have the form

$$g^a_{4i} \approx \frac{\partial A^a_i}{\partial \rho}\dot{\rho} = \frac{4}{g}\bar{\eta}_{ai\nu}\ \frac{y_\nu\ \rho}{(y^2+\rho^2)^2}\ \dot{\rho},\ \ g^a_{ij} = 0\ ,\ \ g^a_{i4} = -g^a_{4i}\ ,\ \ i,j = 1,2,3\ .$$

Here we restrict ourselves with the superposition ansatz (1) and neglect the terms of order $O(\dot{\rho}^2)$. The adiabatic constraint $g^a_{\mu\nu} \ll G^a_{\mu\nu}$ means the variation of instanton size is much smaller than characteristic transformation scale of the PP field, $\dot{\rho} \ll O(1)$. Then calculating the corrections for the action, it is reasonable to take out $\dot{\rho}$ beyond the integral and one instanton contribution to the action turns out

$$s = \frac{1}{4}\int d^4x\ G'^2_{\mu\nu} \simeq \frac{8\pi^2}{g^2} + C\ \dot{\rho} + \frac{\kappa_{s.t.}(\rho)}{2}\ \dot{\rho}^2, \tag{7}$$

where $\dot{\rho}$ should be taken as the mean rate of slow solution deformation at characteristic instanton lifetime $\sim\ \rho$. For simplicity, one may take it in the centre of instanton $\dot{\rho}(z)$. The constant C coming from the interference term should be equal to zero as a result of space homogeneity and for the 'kinematical' κ-term we have [4]

$$\frac{\kappa_{s.t.}(\rho)}{2} = \frac{1}{2}\int d^4x\ g^2_{4i} = \frac{3\pi^2}{g^2}\ . \tag{8}$$

The overt ρ dependence of κ is lacking because of the scale invariance. It arises with the renormalization of the coupling constant. Dealing with definition (1) we have considered only the corrections induced by the variation of strength tensors in line with anzatz (6), but not those

[4] In a regular gauge the result is the same.

resulting from a possible variation of fields (2). Bearing in mind the form of potentials in regular $(r.g.)$ and singular $(s.g.)$ gauges

$$\begin{aligned} A^a_\mu &= \frac{1}{g}\eta_{a\mu\nu}\,\partial_\nu \ln(y^2+\rho^2), \qquad \text{r.g.} \\ A^a_\mu &= -\frac{1}{g}\bar\eta_{a\mu\nu}\,\partial_\nu \ln\left(1+\frac{\rho^2}{y^2}\right), \qquad \text{s.g.} \end{aligned} \tag{9}$$

we find the adiabatic corrections $A'_\mu = A_\mu + a_\mu$ as follows:

$$a^a_\mu = \frac{2}{g}\eta_{a\mu 4}\,\frac{\rho}{y^2+\rho^2}\,\dot\rho, \qquad \text{r.g.} \tag{10}$$

as for the profiles (9), the substitution $\eta \to -\bar\eta$ brings about the transition from regular gauge to a singular one. Consequently, it becomes possible to calculate the correction for the chromoelectric field (in the singular gauge, for clarity)

$$\epsilon^a_i(p.) = g^a_{4i} = \frac{4}{g}\,\frac{\rho\, y_4\, \dot\rho}{(y^2+\rho^2)^2}\,\delta_{ai}\ , \tag{11}$$

(p. means the potential part here), and the correction for the chromomagnetic field is

$$h^a_i = \frac{1}{2}\varepsilon_{ijk}\, g^a_{jk} = -\frac{4}{g}\,\frac{\rho\,\dot\rho}{(y^2+\rho^2)^2}\,\varepsilon_{aij}\, y_j\ . \tag{12}$$

Chromomagnetic field contributes to the action as

$$\frac{1}{2}\int d^4x\; h^a_i h^a_i = \frac{4\pi^2}{g^2}\,\dot\rho^2, \tag{13}$$

and the contribution of chromoelectric field ($\epsilon^a_i = \epsilon^a_i(p.) + \epsilon^a_i(s.t.)$) is given in the following form

$$\frac{1}{2}\int d^4x\; \epsilon^a_i \epsilon^a_i = \frac{12\pi^2}{g^2}\,\dot\rho^2\ . \tag{14}$$

Finally, summing up both contributions to the kinematical κ-term results in

$$\kappa = \frac{32\pi^2}{g^2}\ . \tag{15}$$

The comparison of Eq. (13) and Eq. (14) allows us to conclude that the chromoelectric fields are influenced mainly by dilatational deformations.

The formulae obtained characterize the deformed instanton, and one can see the adiabatic variation of the PP solution does not change the action considerably. A small contribution of order $O(\dot\rho^2)$ only is added. Turning to the $\bar I I$ liquid, let us rewrite the PP ensemble formulae in the compact form of a density functional (as in statistical mechanics). Assuming the number of PP in an ensemble is still appropriate to consider them separately, one may rewrite the superposition Eq. (1) in the form of double summing as

$$A_\mu(x) = \sum_{i=1}^{K}\ \sum_{j=1}^{\triangle N(\rho_i)} A_\mu(x; i, \gamma_j)\ , \tag{16}$$

Here $\triangle N(\rho_i)$ is the PP number of the size ρ in the interval $(\rho_i,\ \rho_i+\triangle\rho)$, K is the number of partitions within $(\rho_{in},\ \rho_{fn})$ and $A_\mu(x;i,\gamma_j)$ is the (anti)-instanton solution for PP with the calibrated size of $\rho\in(\rho_i,\ \rho_i+\triangle\rho)$ where $\gamma=(z,\ \Omega)$ stands for the coordinate of centre and colour orientation of solution. By definition $\sum_{i=1}^{K}\triangle N(\rho_i)=N$. The distribution function is taken as

$$n(\rho)=\frac{\triangle N(\rho)}{\triangle\rho}\frac{1}{V}, \tag{17}$$

with normalizing $\sum_{i=1}^{K} n(\rho_i)\ \triangle\rho\ V=N$. Assuming the limit $\triangle\rho\to 0$ existing one has $V\int d\rho\ n(\rho)=N$. As is known, such an ensemble is described by the action averaged over the PP coordinates in metric and colour spaces[5]. Then it is easy to obtain (for noninteracting system)

$$\langle S^{(0)}\rangle=\prod_{i=1}^{N}\int\frac{d^4z_i}{V}d\Omega_i\ S^{(0)}=\int d^4x\int d\rho\ n(\rho)\left\{\frac{8\pi^2}{g^2}+\frac{\kappa(\bar\rho)}{2}\ \dot\rho^2\right\}\ , \tag{18}$$

where the integration over d^4z is carried out in the volume V occupied by this system, $d\Omega$ is the measure in colour space with unit normalization (we omit the argument ρ in the rate $\dot\rho$ hereafter, hoping it is not misleading). This form of $\langle S\rangle$ reflects the homogeneous property of excitation wave function assumed, together with taking into account possible adiabatic variation of the solution profile in various space points. According to Eq. (18) the deformations at a constant rate additionally favour non-interacting PP's.

The situation becomes crucially different if the interaction is switched on. The pair interaction (and the loop corrections) is expected to be of main importance for our purposes. It is convenient to set the intensity of gluon field $G_{\mu\nu}$ generated by any pair of PP's as a sum of field intensity generated by each separate PP and a term stipulated by nonlinearity of strength tensors [3]

$$\begin{aligned}G_{\mu\nu}(a+b)&=G_{\mu\nu}(a)+G_{\mu\nu}(b)+\triangle G_{\mu\nu}(a,b)\ ,\\ \triangle G_{\mu\nu}(a,b)&=-i\{[a_\mu,b_\nu]+[b_\mu,a_\nu]\}\ ,\end{aligned} \tag{19}$$

(in short $G(1+2)=G_1+G_2+G_{12},\ G_{\mu\nu}=G^a_{\mu\nu}T^a$ and T^a are the group generators). The cumulative field of the ensemble is given by summing up the integral (18) and integrate over all interacting pairs of PP's

$$\begin{aligned}\langle S_{\rm int}\rangle\ &=\ \int d^4x\int\frac{d^4z_1}{V}d\Omega_1\int\frac{d^4z_2}{V}d\Omega_2\\ &\quad\ \frac{1}{2}\sum_{12}\frac{1}{4}\{G_{12}^2+2G_1G_2+2G_1G_{12}+2G_2G_{12}\}.\end{aligned}$$

Colour averaging preserves the first term only and with the explicit form of solution (2) we obtain for the partial contribution of every pair [3]

$$\frac{1}{4}\langle G_{12}^2\rangle=\frac{8\pi^2}{g^2}\frac{N_c}{N_c^2-1}\int d^4x\int\frac{d^4z_1}{V}\int\frac{d^4z_2}{V}\,\frac{[7y_1^2y_2^2-(y_1y_2)^2]\ \rho_1^4\rho_2^4}{y_1^4(y_1^2+\rho_1^2)^2\ y_2^4(y_2^2+\rho_2^2)^2}\ , \tag{20}$$

---

[5] We omit the integration over the instanton sizes because we have introduced the collective coordinate of density. In principle, the generating functional of the theory should be rewritten in terms of functional integration over $D[n(\rho)]$, however, we use here the distribution function (3) found formerly within one-instanton approximation.

where $y_i = x - z_i$, $\rho_i = \rho_i(z_i)$, $i = 1, 2$. In the adiabatic approximation the integral (20) is estimated to be (the contact interaction)[6]

$$\frac{1}{4}\langle G_{12}^2\rangle \simeq \frac{8\pi^2}{g^2}\frac{\xi^2}{V}\int \frac{d^4z}{V}\,\rho_1^2(z)\,\rho_2^2(z)\ .$$

Indeed, the adiabatic condition makes it possible to rescale the integration variable $\frac{dz}{\rho} = d\left(\frac{z}{\rho}\right) + \frac{z}{\rho^2}d\rho \approx d\left(\frac{z}{\rho}\right)$. Besides, it is easy to see that the integral is dominated by the local vicinity of diagonal $\frac{z_1}{\rho_1} \approx \frac{z_2}{\rho_2}$ and the distribution function $n(\rho)$ reduces the interaction term $S_{\rm int}$ to the following form

$$\langle S_{\rm int}\rangle = \frac{8\pi^2}{g^2}\,\xi^2\int d^4z\int d\rho_1\int d\rho_2\,\rho_1^2\,n(\rho_1)\ \ \rho_2^2\,n(\rho_2)\ , \tag{21}$$

remember now ρ_i $(i = 1, 2)$ are the functions of z.

As known, the loop corrections with pre-exponential factor modify the contribution, (18) which is estimated in our approach as (in terms of $\beta = \frac{8\pi^2}{g^2}$ as defined by Eq. (5))

$$\langle S^{(0)}\rangle = \int d^4z\int d\rho\,n(\rho)\left(\frac{\kappa(\bar\rho)}{2}\,\dot\rho^2 + \beta(\rho) + 5\ln(\Lambda\rho) - \ln\widetilde{\beta}^{2N_c}\right)\ . \tag{22}$$

Finally, one has for the averaged action

$$\langle S\rangle = \langle S^{(0)}\rangle + \beta\xi^2\int d^4z\int d\rho\,n(\rho)\,\rho^2\ \int d\rho_2\,n(\rho_2)\,\rho_2^2\ . \tag{23}$$

With the distribution function (3) for stationary ρ we easily reproduce the result (4) taking a conventional quantity of averaged squared instanton size $\overline{\rho^2} = \int d\rho\,\rho^2\,n(\rho)/\int d\rho\,n(\rho)$. Thus, the averaged action is modified to give an effective potential in the following form

$$\langle S\rangle = \int d^4z\int d\rho\,n(\rho)\ \left(\frac{\kappa(\bar\rho)}{2}\,\dot\rho^2 + U_{\rm eff}(\rho)\right)\ ,$$
$$U_{\rm eff}(\rho) = \beta(\rho) + 5\ln(\Lambda\rho) - \ln\widetilde{\beta}^{2N_c} + \nu\,\frac{\rho^2}{\overline{\rho^2}}\ .$$

This result is also used as a suitable approximation herein for the PP's with deformations if it includes an additional integration over volume in the definition of $\bar\rho$ averaged, i.e.

$$\overline{\rho^2} = \frac{\int d^4z\int d\rho\,\rho^2\,n(\rho)}{\int d^4z\int d\rho\,n(\rho)}\ . \tag{24}$$

Everything considered so far concerns the Euclidean action only. In order to get a guess for the Minkowski space we notice that in every temporary slice ($x_4\ =\ const$) we have a dynamical system (with a natural transition from $\dot\rho^2$ in Euclidean space to $-\dot\rho^2$ in Minkowski one) where, as easily seen, appreciable benefit in action could occur. In fact, it takes place if the

[6] If the PP sizes do not vary we have for the partial contribution the known answer [3] (Nucl. Phys.) $\langle S_{\rm int}\rangle = \frac{8\pi^2}{g^2}\ \frac{\xi^2}{V}\ \rho_1^2\rho_2^2$.

rates of solution deformations are self-consistent in every temporary slice and if the oscillator fluctuations of excitations are permissible (just the case of interacting potential above). We imply no fluctuations of instanton solution itself (their lifetime is of instanton size order). Let's emphasize that it is worth speaking about the coherent behaviour of vacuum wave function when the probabilities of appearance and disappearance of instantons with different rate of size deformation are concordant. Besides, the ensemble projection upon the three-dimensional space is described by the Lagrangian density for every group of the PP's of calibrated sizes (see Eq. (16) for partitioning at the initial time)

$$\mathcal{L} = \frac{1}{2}\,\kappa(\bar\rho)\,\dot\rho^2 - U_{\rm eff}(\rho)\ . \tag{25}$$

The characteristic scale of excitations is determined at the point of potential minimum $\frac{dU_{\rm eff}(\rho)}{d\rho} = 0$ and is derived from the equation $\rho_c^2 = \frac{b-5}{2\nu}\,\overline{\rho^2}$. Then $\frac{d^2U_{\rm eff}(\rho)}{d\rho^2}\mid_{\rho_c} = \frac{4\nu}{\overline{\rho^2}}$. Fixing the kinetic energy term (15) on the characteristic scale $\bar\rho$ (similar to the calculation of potential) one has for the frequency of the proper system fluctuations

$$m^2 = \frac{4\nu}{\kappa\,\overline{\rho^2}} = \frac{g^2}{8\pi^2}\frac{\nu}{\overline{\rho^2}} = \frac{\nu}{\beta(\bar\rho)\,\overline{\rho^2}}\ . \tag{26}$$

We have only analyzed the deformations in the temporary direction. Those in the spatial directions could be estimated by drawing the same arguments. Thus, the expression for the κ-term keeps the form above obtained with the only change of rates for the appropriate gradients of function $\rho(x)$ done, i.e. the substitution $\dot\rho(t) \to \frac{\partial\rho(x)}{\partial x}$ should be performed for such a 'crumpled' instanton. Then the frequency of proper fluctuations might be interpreted as the mass term and excitations occur having a phonon-like nature

$$\mathcal{L} = \frac{1}{2}\,\kappa(\bar\rho)\,[\dot\rho^2 - \nabla\rho\nabla\rho] - U_{\rm eff}(\rho)\ , \tag{27}$$

(cross-terms $\sim\ \dot\rho\ \rho^{'}$ equal to zero identically)[7]. The parameters $\bar\rho$ and $\beta(\bar\rho)$ determined (by maximizing the partition function of $\bar I I$ liquid with respect to N [3]) in a self-consistent way take the following values $\bar\rho\Lambda \approx 0.32,\ \ \beta(\bar\rho) \approx 18.1,\ n\ \Lambda^4 \approx 0.73,\ \ (N_c = 3)$ therefore, for mass term we have $m \approx 1.37\Lambda$ [8].

It is a remarkable fact that instanton liquid model gives the effective Lagrangian, incorporating low energy excitations which are naturally interpreted as particles. Their quantum numbers are the scalar 'glueball' ones and the mass gap is of order $\sim\ \Lambda_{\rm QCD}$. This value of the mass scale looks very reasonable and absolutely selfconsistent for an adiabatic approximation within instanton picture of quenched QCD vacuum. Indeed, it is qualitatively clear that 'hard' glueball masses $\sim\ 1.5$ GeV (what is just known from existing theoretical speculations in lattice QCD [5] and sum rule method [6]) wittingly leads beyond our constraints imposed. It means that the strength of deformation field becomes quite compatible with the instanton field and suppresses

[7] It is interesting to remark the centre of instanton solution may not be shifted unlike the dilatational mode, since the relevant deformations make κ singular.

[8] Taking the constant $C_{N_c} \approx \frac{4.66\ \exp(-1.68N_c)}{\pi^2(N_c-1)!(N_c-2)!}$ as in Ref. [3] (Nucl. Phys.) we have $\bar\rho\ \Lambda \approx 0.37,\ \ \beta(\bar\rho) \approx 17.5,\ n\ \Lambda^4 \approx 0.44$, and for the mass gap $m \approx 1.21\ \Lambda$.

the PPs. Unfortunately, nowadays lattice simulation capabilities does not allow us to extract the 'soft' component reliably. The recent comparative analysis of lattice [7] and instanton approach results for correlators [8] supposes to fix only the instanton contribution there available with phenomenologically suitable magnitudes of average instanton size $\bar{\rho}$ and instanton liquid density being calculated in the first order of n-expansion (in fact, packing fraction-expansion $n\ \rho^4$). But already the next order of this expansion is useless because of impossibility to reproduce even the proper signs of corresponding correlation function components. The main difficulty at picking out of it comes from the presence of huge perturbative background together with lack of reliable theoretical estimates in the long distance region. Besides, the direct use of standard lattice simulations to isolate the $O(\dot{\rho}^2)$ (crumpled) contribution requires the measurements at now inconceivably large values of $\beta \sim$ 12–18 and the cooling technique is 'ironing the crumpling'. However, our analysis predicts the interesting feature in the behaviour of crumpled correlators. Their chromoelectric components occur three times larger than the chromomagnetic ones [9].

Let's emphasize once more that adiabatic assumption leads to fully consistent picture of $\bar{I}I$ liquid. The model itself regulates the most suitable regime of dilatation deformations resulting in the mass gap generation (26). Apparently, an intriguing guess is to associate some class of light hadrons with these phonon excitations discovered, since the preliminary evaluations of their mass spectrum with the quark condensate included look quite encouraging. The concept of confining potential origin for light quarks in this approach seems simply irrelevant because of stable phonon nature.

Acknowledgement The Financial support of RFFI: Grants 96-02-16303, 96-02-00088 G, 97-02-17491 and INTAS Grant 93-0283, 96-0678 is greatly acknowledged.

References

[1] M. Alford, K. Rajagopal, F. Wilczek: *hep-ph/9711395*, *Phys. Lett.* **B422** (1998) 247;
R.Rapp, T. Schäfer, E. V. Shuryak, M. Velkovsky: *hep-ph/9711396*, *Phys. Rev. Lett.* (in press).

[2] C. G. Callan, R. Dashen, D. J. Gross: *Phys. Rev.* **D17** (1978) 2717;
C. G. Callan, R. Dashen, D. J. Gross: *Phys. Lett.* **B66** (1977) 375;
E.-M. Ilgenfritz, M. Müller-Preussker: *Nucl. Phys.* **B184** (1981) 443;
E. V. Shuryak: *Nucl. Phys.* **B203** (192) 93, 116, 140; **B328** (1989) 85, 102;
T. Schäfer, E. V. Shuryak: *Rev. Mod. Phys.* **70** (1998) 323

[3] D. I. Diakonov, V. Yu. Petrov: *Nucl. Phys.* **B245** (1984) 259;
D. I. Diakonov, V. Yu. Petrov: in: *Hadron matter under extreme conditions*, ed. by V. Shelest and G. Zinovjev, Naukova Dumka, Kiev (1986) p. 192;
D. I. Diakonov, V. Yu. Petrov, P. V. Pobylitsa: *Phys. Lett.* **B226** (1989) 471

[4] G. 't Hooft: *Phys.Rev.* **D14** (1976) 3432

[5] M. Teper: *hep-ph/9711299*
W. Lee, D. Weingarten: *hep-lat/9805029*

[6] Tao Huang, Hong Ying Jin, Ailin Zhang: *hep-ph/9807391*

[9] At finite temperature they linearly raise with temperature increasing, if the kinetic term κ is constant, and the caloron, which is temperature dependent itself, should be substituted in lieu of instanton (it is in progress now).

[7] A. Di Giacomo, H. Panagopoulos: *Phys. Lett.* **B285** (1992) 133;
L. Del Debbio, A. Di Giacomo, Yu. A. Simonov: *Phys. Lett.* **B332** (1994) 111;
A. Di Giacomo, E. Meggiolaro, H. Panagopoulos: *hep-lat/9603017*;
Talk given at *High-energy Physics International Euroconference on Quantum Chromodynamics* (QCD 96), Montpellier, France, 4-12 Jul 1996, *Nucl. Phys. Proc. Suppl.* **54A** (1997) 343;
A. Di Giacomo, E. Meggiolaro, H. Panagopoulos: *Nucl. Phys.* **B483** (1997) 371;
M. D'Elia, A. Di Giacomo, E. Meggiolaro: *Phys. Lett.* **B408** (1997) 315

[8] E.-M. Ilgenfritz, B.V. Martemyanov, S.V. Molodtsov, M. Müller-Preussker, Yu.A. Simonov: *hep-ph/9712523*, *Phys. Rev.* **D58** (1998) (id number 114508).