1. INTRODUCTION

In spite of the tremendous experimental and theoretical efforts in the investigation of the properties of quark-gluon plasma (QGP) at large temperatures and densities, it is not clear so far what is the fundamental QCD mechanism leading to the unusual behaviour of the matter produced in heavy ion collisions at high energies at the RHIC and LHC. Indeed, there is strong evidence that even above $T_c$ non-perturbative QCD effects are very important and, in particular, they are responsible for the phenomenon of the so-called strongly-interacting Quark-Gluon Plasma (sQGP) [1] observed at the RHIC and LHC. One of the important outputs of the Lattice QCD calculation at finite $T$ is the fact that the values of the deconfinement and chiral restoration temperatures are approximately equal. That means that above $T_c$ one cannot expect the pion to be the Goldstone boson with zero mass in the chiral limit. Instead, a rather massive ($M_\pi \approx 2M_q(T_c) \approx 6T_c \sim 1$ GeV) pion state appears above $T_c \approx 150$ MeV. Therefore, this lowest mass quark-antiquark state can not give a significant contribution to the Equation of State (EoS) of the QGP.

One of the fundamental issues of QCD as the theory of strong interactions is the understanding of the role of gluonic degrees of freedom in the confinement and deconfinement regimes. Thus, it is known that at low temperatures gluons play an important role not only in the dynamics of usual hadrons. In particular, they can form bound states called glueballs (see review [2]). However, one cannot expect a significant contribution of glueballs at $T < T_c$ to EoS of the hadron gas due to their large masses $M_G \gg T$. On the other hand, it was found by lattice calculation [3, 4] that unlike the quark condensate, the gluon condensate does not vanish at $T > T_c$. We should stress that this condensate plays a fundamental role in the formation of the bound glueball states. This role is similar to the role of the quark condensate in the appearance of the massive mesons and baryons built of the light $u$, $d$, and $s$ quarks. Therefore, a non-zero value of the gluon condensate above $T_c$ is a strong signal of the existence of the glueballs in the deconfinement phase.

However, it is evident that the properties of the glueballs, in particular, their masses and sizes should change in the QGP due to the temperature dependence of the gluon condensate [4, 5] and the change of the topological structure of the QCD vacuum at $T > T_c$ [5, 6]. The attempts to estimate the value of the glueball masses at finite $T$ were done in the different models in the papers [16–20] and in the lattice [21, 22].

In this Letter, we consider the lowest scalar and pseudoscalar glueball states at finite $T$ in the QGP environment within the effective model based on the instanton picture for the QCD vacuum. It is shown that their masses strongly decrease above $T_c$ and become very small at the temperature of the scale invariance restoration $T_{scale} \approx 1$ GeV. The possibility of the Bose–Einstein condensation and superfluidity of the glueball matter is under discussion.

1 The article is published in the original.
2. NONPERTURBATIVE QUARK–QUARK, QUARK–GLUON AND GLUON–GLUON INTERACTIONS BELOW AND ABOVE Tc INDUCED BY INSTANTONS

It is expected that the instantons, a strong vacuum fluctuation of gluon fields, play a very important role in the dynamics of the glueballs below Tc (see review [23]). The effective interaction induced by the instanton between quarks at T = 0 is well known. This is a famous t’Hooft interaction which for $N_f = 3$ (Fig. 1a) and $N_c = 3$ is given by the formula [26]:

$$ L_{\text{eff}}^{(3)} = \int d\rho(n(\rho) \prod_{i=a,d,s} \left[ m_i^{\text{car}} \rho - \frac{4\pi^2}{3} \frac{\rho^3}{q_{i\ell A}} \right]$$

$$+ \frac{3}{32} \left( \frac{4\pi^2}{3} \rho^3 \right) \left[ j_{\mu\nu} j_{\mu\nu} - \frac{3}{4} j_{\mu\nu} j_{\mu\nu} \right]$$

$$\times \left( \beta_{\mu\nu} \sigma_{\mu\nu} + \text{perm} \right) + \frac{9}{320} \left( \frac{4\pi^2}{3} \rho^3 \right) d_{abc} j_{\mu\nu} j_{\mu\nu}$$

$$+ \frac{i f_{abc}}{256} \left( \frac{4\pi^2}{3} \rho^3 \right) d_{abc} j_{\mu\nu} j_{\mu\nu} + (R \leftrightarrow L),$$

(1)

where, $m_i^{\text{car}}$ is the quark current mass, $q_{R,L} = (1 \pm \gamma_5)q(x)/2$, $j_{\mu\nu} = \bar{q}_R \lambda^a \gamma_{\mu\nu} q_L$, $j_{\mu\nu} = \bar{q}_R \sigma_{\mu\nu} \lambda^a q_L$, $\rho$ is the instanton size and $n(\rho)$ is the density of instantons. For $N_f = 2$ the t’Hooft interaction for zero current quark mass is much simpler

$$ L_{\text{eff}}^{(N_f=2)} = \int d\rho(n(\rho) \prod_{i=a,d,s} \left[ m_i^{\text{car}} \rho - \frac{4\pi^2}{3} \frac{\rho^3}{q_{i\ell A}} \right]$$

$$\times \left[ 1 + \frac{3}{32} \lambda^a \lambda^a \sigma_{\mu\nu} + \frac{9}{32} \sigma_{\mu\nu} \sigma_{\mu\nu} \lambda^a \lambda^a \right] + (R \leftrightarrow L).$$

(2)

In this section, all Lagrangians are written in the Euclidian space-time.

This Lagrangian can be obtained from Eq. (1) by connecting the strange quark legs through the quark condensate or by the current quark mass.

The effective quark-gluon interaction induced by the instantons is (see, for example [27, 28])

$$ L_{\text{ggqf}} = \int d\rho(n(\rho) \prod_{i=a,d,s} \left[ m_i^{\text{car}} \rho - \frac{4\pi^2}{3} \frac{\rho^3}{q_{i\ell A}} \right]$$

$$\times \left[ (G^2 - G\tilde{G}) L_{f,A} + (G^2 + G\tilde{G}) L_{f,A} \right],$$

(4)

where

$$ G^2 \equiv G_{\mu\nu}^a G_{\mu\nu}^a, \quad G\tilde{G} \equiv G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a,$$

(5)

and

$$ L_{f,A} = (R \leftrightarrow L),$$

(6)

Due to factor $2\pi^2 / g_s^2$ in Eq. (3) each two gluon legs produce a large enhancement factor $\pi^2 / g_s^2$ in the

Fig. 1. The quark–quark (a), quark–gluon (b), (c) and gluon–gluon interaction (d) induced by the instanton for the $N_f = 2$ case. The symbol $\times$ means connection through the quark condensate or by the current quark mass.
gluon-gluon interaction induced by instantons. So one can expect that this type of the interaction is much stronger in comparison with the quark-quark case.

We would like to emphasize that the instanton induced interaction is very sensitive to the parity of the glueball and quarkonium states and, in particular, it leads to the mass splitting between scalar and pseudoscalar glueballs [29, 37]. Indeed, the single instanton contribution to the difference of two correlators of glueball currents with opposite parities is given by

$$\Delta \Pi(Q^2)_{G} = i \int dx e^{i a} \left\langle \langle 0 | T O_{s}(x) O_{s}(0) | 0 \rangle \right\rangle - \left\langle \langle 0 | T O_{p}(x) O_{p}(0) | 0 \rangle \right\rangle = 2 \pi^2 \int d\rho \mathcal{P}(\rho) \mathcal{P}(Q)^2 K_2^2(\rho Q), \tag{7}$$

where the glueball currents for the scalar and pseudoscalar states are the following:

$$O_{s}(x) = \alpha_{s} G_{\mu \nu}^{s}(x) G_{\mu \nu}^{s}(x), \tag{8}$$

$$O_{p}(x) = \alpha_{p} G_{\mu \nu}^{p}(x) G_{\mu \nu}^{p}(x). \tag{9}$$

At \( T < T_c \), the main contribution to the quark-quark and gluon-gluon interaction comes from disordered instantons, Fig. 1, in the so-called “random” phase of the instanton liquid. The phase provides spontaneous chiral symmetry breaking in QCD (see reviews [23, 24]). In this case, the chiral symmetric contribution, which is related to the correlated instanton-antiinstanton molecules, Fig. 2, is expected to be small due to the small packing fraction of instantons in the QCD vacuum [23]. Above \( T_c \), where the chiral and \( U_A(1) \) symmetries are restored, the situation is opposite. Indeed, in this case the contribution from the chirality violated interaction, Fig. 1, is proportional to the product of the current quark masses and should be small. In the chiral symmetric phase, the leading contribution comes from the strongly correlated instanton-antiinstanton molecules, Fig. 2, as supported by the calculations in [25]. The contribution from the eight-quark interaction presented in Fig. 3a does not give a significant contribution to the hadron spectroscopy but it is important for the stability of the vacuum of the NJL model with the instanton induced interaction [30, 31].

The effects of quark-quark interaction induced by instanton-antiinstanton molecules near \( T_c \) were studied in [32, 33]. It was shown that they might give a significant contribution to the binding energy of the quark-antiquark states in the deconfinement phase. The main effect, in comparison with the zero temperature case, comes from the strong polarization of instanton-antiinstanton molecules in the time direction at the high temperature.

3. LOWEST MASS SCALAR GLUEBALL ABOVE \( T_c \)

At \( T = 0 \) there are several scalar meson states \( f_{0}(600), f_{0}(980), f_{0}(1500) \) and \( f_{0}(1710) \) with a possible
To estimate the coupling $\lambda$, we use the relation
$$\lambda = \frac{4\pi n_0^4}{9\langle \alpha_s G^2 \rangle},$$
and the value of the gluon condensate [14]
$$\langle \alpha_s G^2 \rangle \approx 0.07 \text{ GeV}^4.$$ 

For the mass of the glueball at $T \approx T_c$ $m_0 \approx 450$ MeV we have $\lambda = 0.82$. Now we are in a position to calculate the temperature dependence of the scalar glueball mass above $T_c$. At the finite temperature, the effective potential in the one loop approximation for the $\lambda \Phi^4$ theory is given (see for example, [15])
$$V(\Phi(T)) = V_0(\Phi(T)) + \frac{3\lambda \Phi(T)}{4\pi^2} \int \frac{k^2 dk}{\sqrt{k^2 + m_0^2} \exp(\sqrt{k^2 + m_0^2 / T} - 1)}.$$ 

By using the condition of the minimum of $V(\Phi(T))$
$$\frac{\delta V(\Phi(T))}{\delta \Phi(T)} = 0,$$
and the definition for the mass of the glueball
$$\frac{\delta^2 V(\Phi(T))}{\delta \Phi(T)^2} = m_0^2(T),$$
we obtain
$$m_0^2(T) = m_0^2 - \frac{3\lambda}{\pi^2} \int \frac{k^2 dk}{\sqrt{k^2 + m_0^2} \exp(\sqrt{k^2 + m_0^2 / T} - 1)}.$$ 

For the trace anomaly we obtain
$$T^\mu_\mu = 4V(\Phi_{\text{min}}(T)) - \lambda \Phi_{\text{min}}(T)^4 = -\frac{m_0^4(T)}{4\lambda}.$$ 

The temperature dependence of the gluon condensate in this model is given by the formula
$$\langle \alpha_s G^2(T) \rangle = \frac{4m_0^4(T)}{9\lambda}.$$ 

In Fig. 4, the temperature dependence of the mass of the scalar glueball is presented. So one can see that at $T_{\text{scale}} \approx 0.9 \text{ GeV}$ the mass of the scalar glueball vanishes and the scale invariance is restored. This can be treated as the appearance of the massless dilaton field in the limit $T^\mu_\mu \to 0$ (see the recent discussion of the dilaton in [43]).

We should emphasize that the justification of the use of the point-like effective interaction in Eq. (11) can be related to a very small size of the scalar glueball.
in the QGP. Even at $T = 0$, due to the strong attraction between gluons induced by the instantons, the size of the scalar glueball is very small, $R_\text{c} \approx 2/3\rho_\text{c} \approx 0.2$ fm, where $\rho_\text{c} \approx 0.3$ fm is the average instanton size in the QCD vacuum [34]. This small size was also confirmed in the lattice calculation [35, 36]. Above $T_\text{c}$, a similar phenomenon happens as well. However, in this case the instanton-antiinstanton molecules produce a very strong attraction in the scalar glueball channel. Therefore, we expect that the size of the scalar glueball should be smaller than the size of instantons in the QGP, which is cutting at

$$R_\text{c} \approx p^2(T) \approx \frac{1}{3\pi^2 T^2}, \quad (19)$$

at finite T and $N_c = N_f = 3$ [39]. We would like to emphasize that this size is smaller than the perturbative Debye screening length in the QGP

$$\lambda^2(T) = \frac{1}{M^2(T)} > \frac{1}{3\pi T^2}, \quad (20)$$

where $a_0(\rho) \approx 0.5$ [24] for $T = 0$ and $M^2(T) = g^2 \nu (N_c/3 + N_f / 6)T^2$ [38] was used. Therefore, we come to the important conclusion that the Debye screening cannot destroy the binding of the scalar glueball in QGP.

4. PSEUDOSCALAR GLUEBALL ABOVE $T_\text{c}$

There are also several candidates for the pseudoscalar glueball below 2 GeV at $T = 0$: $\eta(1405)$, $\eta(1760)$, $\eta(1475)$, and $X(1835)$. In the lattice quenched calculation the lowest mass pseudoscalar glueball was predicted to have the mass $M_0^2 \approx 2.6$ GeV [40, 41]4. The large difference between scalar and pseudoscalar glueball masses $m_0^2(T = 0) - m_0^2(T = 0) \approx 1$ GeV at low temperature can be explained by the large single instanton contribution to these channels (see discussion in [23]). Indeed, it gives a strong attraction in the scalar channel and a strong repulsion in the pseudoscalar state. After restoration of the chiral symmetry the masses of scalar and pseudoscalar glueballs should be equal in the limit of the zero light quark masses, $m_u = m_d = m_s = 0$. In this case, instanton-antiinstanton molecules give the same strength of attraction in the both states. Therefore, the leading contribution to the splitting between the masses of two states at $T > T_\text{c}$ is determined by the density of the single instantons, which is proportional to the product of the current masses of the light quarks. We should mention that even at large temperature the single instanton contribution is repulsive in the pseudoscalar glueball case, and its collapse to the massless state is not allowed. This is a very important difference between the properties of pseudoscalar and scalar glueballs at the temperature of the scale invariance restoration.

We can estimate the mass of the pseudoscalar glueball at $T = T_{\text{scale}}$ by using the instanton model. The difference between the scalar and pseudoscalar glueball masses above $T_\text{c}$ is determined by the instanton density $n(\rho(T))$

$$n(\rho(T)) = m_\text{sc} m_\text{c} (\rho(T))^{b_0 - 2}, \quad (21)$$

where $b_0 = 11N_c/3 - 2N_f / 3$ and $m_\text{c}$ are the current masses of the quarks. Finally, we have

$$m_0^2(T_{\text{scale}}) = \frac{m_0^2 m_\text{sc} m_\text{c}}{m_\text{sc}^2 m_\text{c}^2 m_s^2} \left(\frac{1}{3\pi T_{\text{scale}} \rho_\text{c}}\right)^7 \times (m_0^2(T = 0) - m_0^2(T = 0)), \quad (22)$$

$$m_\text{c} = m_u + m_d$$

where $m_\text{c}$ is the so-called effective mass of the quark in the instanton vacuum related to the quark condensate. The use of $m_u = m_d \approx 4.5$ MeV, $m_s \approx 100$ MeV [42], $m^* = 170$ MeV [23], $\rho_\text{c} = 1/600$ MeV$^{-1}$ and $m_\text{c} = m_0^2(T = 0) - m_0^2(T = 0) \approx 1$ GeV gives the estimation

$$m_\text{c} \approx 0.1 \text{eV}. \quad (23)$$

Therefore, the mass of the pseudoscalar glueball at $T = T_{\text{scale}}$ is finite but very small.

5. BOSE–EINSTEIN CONDENSATION OF THE LIGHT GLUEBALLS AND SUPERFLUIDITY OF GLUEBALL MATTER IN QGP

The Bose–Einstein condensation (BEC) of the identical bosons is the well-known phenomenon in both theoretical and experimental physics. Recently, it was suggested that the over-occupied initial state of gluons created during relativistic heavy ion collisions might lead to the BEC of gluons [44–47]. However, gluons have the color charge and spin and, additionally,
the number of gluons is not conserved. All of these features, lead to a very complicated study of the possible formation of the gluon BEC in the QGP. On the other hand, the light glueballs with zero spin and color can easily form BEC at rather high temperature because for the given boson number density $n$ the BEC temperature depends on the mass as [48]

$$T_{BEC} \approx 3.31 \frac{n^{3/3}}{n},$$

(24)

for the non-relativistic case. For the arbitrary boson mass the critical BEC density is related to the temperature by the equation (see, for example, [49] and references therein)

$$n_{BEC} = \frac{\int d^3k}{(2\pi)^3} \exp \left( \frac{1}{\sqrt{k^2 + m^2} - m} / T_{BEC} \right) - 1$$

(25)

In Fig. 5, we show the critical glueball density of the scalar and pseudoscalar glueballs in the QGP by using the T-dependent glueball mass from Eq. (16)$^5$. This density can be compared with the density of QGP at the time $\tau$ produced at the RHIC, which was estimated by using the multiplicity of the meson production in the Au-Au central collision [50]

$$n_{RHC} \approx \frac{13}{\tau} (fm^{-3}),$$

(26)

where $\tau$ is in fm. For LHC Pb–Pb central collisions at $\sqrt{s_{NN}} = 2.76$ GeV the multiplicity is approximately twice larger. Therefore, the density is

$$n_{LHC} \approx \frac{26}{\tau} (fm^{-3}).$$

(27)

For QGP in the thermal equilibrium we can estimate the contribution of the glueballs to the total density. In this case, we have approximately 16 gluon and 24 quark degrees of freedom ($N_f = 2$) and 2 for glueballs. For the gluons and quarks with the mass $M_{gg} \approx 3T$ [51], we have for scalar (pseudoscalar) glueball contribution to the total density at the RHIC

$$n_{g^\text{glueball}}^{\text{RHIC}} \approx \frac{3}{\tau} (fm^{-3}),$$

(28)

for the initial temperature $T_0^{\text{RHIC}} \approx 300$ MeV.

For the LHC energy ($T_0^{\text{LHC}} \approx 400$ MeV) the estimation is

$$n_{g^\text{glueball}}^{\text{LHC}} \approx \frac{5}{\tau} (fm^{-3}).$$

(29)

By using the Bjorken model for the QGP expansion [52]

$$\tau T^3 = \tau_0 T_0^3,$$

(30)

with the thermalization time $\tau_0 \approx 1$ fm, we can estimate BEC temperature for the RHIC and LHC

$$T_{BEC}^{\text{RHIC}} \approx 200 \text{ MeV}, \ T_{BEC}^{\text{LHC}} \approx 270 \text{ MeV}.$$  

(31)

Therefore, the formation of the glueball BEC is possible at both RHIC and LHC heavy ion collisions. We should also mention that there is a large gap between the mass of the light glueball and the mass of the excited state in the system of two gluons in the QGP, which is $m_{gg} \approx 6T > 1$ GeV. It is well know that such a gap should lead to the phenomenon of the superfluidity of the BEC matter. So we arrive at the conclusion that the QGP at the RHIC and LHC might be considered as the mixture of three matters. One of them is the usual “normal” QGP matter consisting of quarks and gluons and other ones are two superfluid matters of very light scalar and pseudoscalar glueballs.

There is also an additional mechanism which can lead to the abundance of the light glueballs production in ultra-relativistic heavy ion collisions. Indeed, at very high temperatures at the initial stage of the production of the quark-gluon matter, the pair production of glueballs is possible by the fusion of two gluons, Fig. 6. The effective interaction responsible for such production is

$$\mathcal{L} = \lambda_2(T) \alpha_s G_{\mu\nu}^a G^{\mu\nu}_a (S^2 + P^2),$$

(32)
where $S(\rho)$ are the scalar (pseudoscalar) fields. In the mean field approximation we have

$$\lambda_S(\rho) \sim \frac{m_{S,\rho}(\rho)^2}{\alpha G_{\mu\nu}^a G_{\mu\nu}^a(\rho)} \sim \frac{1}{m_{S,\rho}(\rho)}, \quad (33)$$

Therefore, one might expect a strong enhancement of the ultra-light glueball production in the phase where the quark-gluon plasma is far away from the equilibrium.

6. CONCLUSIONS

In summary, we considered the properties of scalar and pseudoscalar glueballs in the Quark-Gluon Plasma created in relativistic heavy-ion collisions. Based on the instanton model for the QCD vacuum we gave the arguments in favor of the existence of very light scalar and pseudoscalar glueball states above the temperature of the deconfinement transition. The estimation of the temperature of the scale invariance restoration, at which the scalar glueball becomes massless, was given. We also discussed the mechanism of the Bose–Einstein condensation and the superfluidity of the scalar and pseudoscalar glueball matter in the QGP. We also showed the possibility of the abundance glueball production at the initial stage of the QGP formation. The influence of this phenomenon on the fast thermalization of the QGP is the subject of our future investigation.

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