Physics Letters B 746 (2015) 311–314

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Turbulent meson condensation in quark deconfinement

Koji Hashimoto [a,b], Shunichiro Kinoshita [c], Keiju Murata [d], Takashi Oka [e]

[a] Department of Physics, Osaka University, Toyonaka, Osaka 560-0043, Japan
[b] Mathematical Physics Lab., RIKEN Nishina Center, Saitama 351-0198, Japan
[c] Osaka City University Advanced Mathematical Institute, Osaka 558-8585, Japan
[d] Keio University, 4-1-1 Hiyoshi, Yokohama 223-8521, Japan
[e] Department of Applied Physics, University of Tokyo, Tokyo 113-8656, Japan

ARTICLE INFO

Article history:
Received 21 January 2015
Received in revised form 2 April 2015
Accepted 1 May 2015
Available online 6 May 2015
Editor: M. Cvetič

ABSTRACT

In a QCD-like strongly coupled gauge theory at large N_c, using the AdS/CFT correspondence, we find that heavy quark deconfinement is accompanied by a coherent condensation of higher meson resonances. This is revealed in non-equilibrium deconfinement transitions triggered by static, as well as quenched electric fields even below the Schwinger limit. There, we observe a "turbulent" energy flow to higher meson modes, which finally results in the quark deconfinement. Our observation is consistent with seeing deconfinement as a condensation of long QCD strings.

1. Introduction

Quark confinement is one of the most fundamental and challenging problems in elementary particle physics, left unsolved. Although quantum chromodynamics (QCD) is the fundamental field theory describing quarks and gluons, their clear understanding is limited to the deconfined phase at high energy or high temperature limits due to the asymptotic freedom. We may benefit from employing a more natural description of the zero temperature hadron vacuum. A dual viewpoint of quark confinement in terms of the "fundamental" degrees of freedom at zero temperature – mesons, is a plausible option.

The mesons appear in families: they are categorized by their spin/flavor quantum numbers, as well as a resonant excitation level n giving a resonance tower such as $\rho(770), \rho(1450), \rho(1700), \rho(1900), \ldots$. In this Letter we find a novel behavior of the higher meson resonances, *i.e.*, mesons with large n. In the confined phase, when the deconfined phase is approached, we observe *condensation of higher mesons*. In this state, macroscopic number of the higher meson resonances, with a characteristic distribution, are excited. The condensed mesons have the same quantum number as the vacuum. The analysis is done via the anti-de Sitter space (AdS)/conformal field theory (CFT) correspondence [1–3], one of the most reliable tools to study strongly-coupled gauge theories. By shifting our viewpoint from quark–gluon to meson degrees of freedom, we gain a simple and universal understanding of the confinement/deconfinement transition, with a bonus of solving mysteries in black holes physics through the AdS/CFT.

The system we study is the $\mathcal{N}=2$ supersymmetric $SU(N_c)$ QCD which allows the simplest AdS/CFT treatment [4]. The deconfinement transition is induced by external electric fields. In static fields, the confined phase becomes unstable in electric fields stronger than the Schwinger limit $E=E_{\rm Sch}$ beyond which quarks are liberated from the confining force. We find that this instability is accompanied by the condensation of higher mesons. A striking feature is revealed for the case of an electric field quench: The kick from the quench triggers a domino-like energy transfer from low to high resonant meson modes. This leads to a dynamical deconfinement transition [5] even below the Schwinger limit. The transfer we find resembles that of turbulence in classical hydrodynamics as higher modes participate; thus we call it a "turbulent meson condensation" and suggest it being responsible for deconfinement.

We remind that the $\mathcal{N}=2$ theory is a toy model: the meson sector is confined and has a discrete spectrum while the gluon sector is conformal and is always deconfined. It resembles heavy quarkonia in a gluon plasma. Generically, quark deconfinement and gluon deconfinement can happen separately, as is known through charmonium experiments in heavy ion collisions. Here, we concentrate on the deconfinement of heavy quarks and not the gluons. Note that the mesons with low spins in this theory are described by a confining potential [6] and an effective QCD string exists, whereby we define our "quark confinement".

E-mail address: koji@phys.sci.osaka-u.ac.jp (K. Hashimoto).

http://dx.doi.org/10.1016/j.physletb.2015.05.004

The higher meson resonances are naturally interpreted as long QCD strings, therefore our finding is consistent with interpreting deconfinement as condensation of QCD strings [7] (see [8–11]). Under the condensation, a quark can propagate away from its partner antiquark by reconnecting the bond QCD string with the background condensed strings. The gravity dual of the deconfined phase is with a black hole, so given the relation with long fundamental strings [9], our result may shed light on the issue of quantum black holes; In particular, our time-dependent analysis gives a singularity formation on the flavor D-brane in AdS, a probe-brane version of the Bizon–Rostworowski turbulent instability in AdS geometries [12].

2. Review: meson effective action from AdS/CFT

The effective field theory of mesons can be obtained for the $\mathcal{N}=2$ supersymmetric QCD in the large N_c, $\lambda \equiv N_c g_{\rm YM}^2$ limits by the AdS/CFT correspondence [6,13]. The meson action is nothing but a D7-brane action in the $AdS_5 \times S^5$ geometry:

$$S = \frac{-1}{(2\pi)^6 g_{\rm YM}^2 l_s^8}\int d^8\xi \sqrt{-\det(g_{ab}[w] + 2\pi l_s^2 F_{ab})}\,,$$
$$ds^2 = \frac{r^2}{R^2}\eta_{\mu\nu}dx^\mu dx^\nu + \frac{R^2}{r^2}\left[d\rho^2 + \rho^2 d\Omega_3^2 + dw^2 + d\bar{w}^2\right], \qquad (1)$$

where $r^2 \equiv \rho^2 + w^2 + \bar{w}^2$, $F_{ab} = \partial_a A_b - \partial_b A_a$, and the AdS_5 curvature radius is $R \equiv (2\lambda)^{1/4} l_s$. For the following calculations, it is convenient to define a rescaled gauge potential $a_a \equiv 2\pi l_s^2 R^{-2} A_a$. The D7-brane worldvolume fields are $w(x^\mu, \rho)$ and $a_a(x^\mu, \rho)$. (We set $\bar{w}(x^\mu, \rho) = 0$ consistently because of $U(1)$-symmetry in $(w, \bar{w})$-plane.) We denote the location of the D7-brane at the asymptotic AdS boundary as $w(x^\mu, \rho = \infty) = R^2 m$. Here, the constant m is related to the quark mass m_q as $m_q = (\lambda/2\pi^2)^{1/2} m$. A static solution of the D7-brane in the $AdS_5 \times S^5$ geometry is given by $w(x^\mu, \rho) = R^2 m$ and $a_a(x^\mu, \rho) = 0$. Using the AdS/CFT dictionary, normalizable fluctuations around the static solutions of w and a_a are interpreted as the infinite towers of scalar mesons $\bar{\psi}\psi$ and vector mesons $\bar{\psi}\gamma_\mu\psi$, respectively. (We omit the pseudo-scalar mesons which are irrelevant to our discussion.) In this paper, we focus on the meson condensation induced by an electric field along the x-direction. Thus, we only consider fluctuations of w and a_x.

To derive the meson effective action, we use a coordinate z defined by $\rho = R^2 m\sqrt{1-z^2}/z$ (where $z = 0$ is the AdS boundary, and $z = 1$ is the D7-brane center that is closest to the Poincaré horizon in the bulk AdS). The worldvolume is effectively in a finite box along the AdS radial direction, to give a confined discrete spectrum as we will see below. We expand the D7-brane action up to second order in the fluctuations $\boldsymbol{\chi} \equiv (R^{-2}w - m, a_x)$ as [6]

$$S = \int dt d^3x \int_0^1 dz \frac{1-z^2}{2z}[\dot{\boldsymbol{\chi}}^2 - m^2(1-z^2)\boldsymbol{\chi}'^2] + \mathcal{O}(\boldsymbol{\chi}^3),$$

where $\dot{} \equiv \partial_t$ and $' \equiv \partial_z$. An irrelevant overall factor is neglected. The equation of motion for $\boldsymbol{\chi}$ is

$$\left(\partial_t^2 + \mathcal{H}\right)\boldsymbol{\chi} = 0\,, \quad \mathcal{H} \equiv -m^2 \frac{z}{1-z^2}\frac{\partial}{\partial z}\frac{(1-z^2)^2}{z}\frac{\partial}{\partial z}\,. \qquad (2)$$

The eigenfunction of $\mathcal{H}$ is given by $e_n(z) \equiv \sqrt{2(2n+3)(n+1)(n+2)}\, z^2 F(n+3, -n, 2; z^2)$, with the eigenvalue $\omega_n^2 = 4(n+1)(n+2)m^2$, for the meson level number $n = 0, 1, 2, \ldots$. Here F is the Gaussian hypergeometric function. The inner product is defined as $(f, g) \equiv \int_0^1 dz\, z^{-1}(1-z^2) f(z) g(z)$, where $(e_n, e_m) = \delta_{mn}$ is satisfied. Note that an external electric field term, $a_x = -Et$,

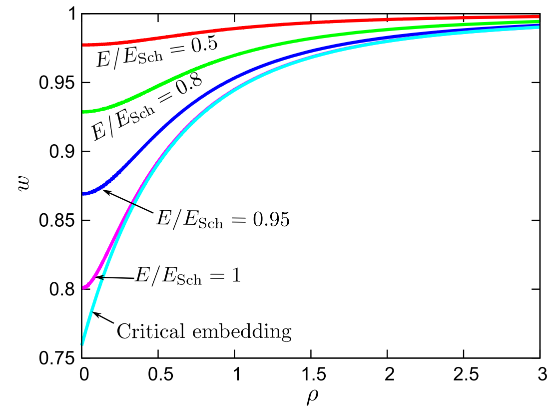

Fig. 1. (Color online.) The shape of the probe D7-brane in static electric fields in the unit of $R = m = 1$. The lines correspond respectively to $E/E_{\rm Sch} = 0.5, 0.8, 0.95, 1$ and the critical embedding from top to bottom.

satisfies Eq. (2) although it is non-normalizable. Expanding the scalar/vector fields as

$$\boldsymbol{\chi} = (0, -Et) + \sum_{n=0}^{\infty} \mathbf{c}_n(t) e_n(z)\,, \qquad (3)$$

we find an infinite tower of meson fields $\mathbf{c}_n(t)$ sharing the same quantum charge – higher meson resonances. Substituting Eq. (3) back to Eq. (1), we obtain the meson effective action

$$S = \frac{1}{2}\int d^4x \sum_{n=0}^{\infty}\left[\dot{\mathbf{c}}_n^2 - \omega_n^2 \mathbf{c}_n^2\right] + \text{interaction}\,, \qquad (4)$$

where we have omitted a constant term and total derivative terms, while higher order nonlinear terms give rise to meson–meson interactions. From the effective action, we obtain the energy $\varepsilon_n \equiv \frac{1}{2}(\dot{\mathbf{c}}_n^2 + \omega_n^2 \mathbf{c}_n^2)$ stored in the n-th meson resonance, and the linearized total energy $\varepsilon = \sum_{n=0}^{\infty} \varepsilon_n$.

3. Higher meson condensation and deconfinement

The confined phase becomes unstable in strong static electric fields. Here, we examine this from the viewpoint of meson condensation. In Eq. (4), the meson couplings depend on the external electric field E nonlinearly. Mesons in a single flavor theory is neutral under E, but it can polarize, and non-linear E may cause a meson condensation. We first solve the equations of motion obtained from the full nonlinear D-brane action (1) with an external electric field, and then decompose the solution $\boldsymbol{\chi}(t, z)$ as Eq. (3). In this way, we can study how the infinite tower of mesons $\mathbf{c}_n(t)$ behave towards the deconfinement transition.[1]

The static D7-brane solution in the presence of a constant electric field introduced by $a_x = -Et$ was obtained in Refs. [14–16]. Also, the Schwinger limit $E = E_{\rm Sch} = 0.5759\ \mathrm{m}^2$ beyond which the first order phase transition to deconfinement occurs was found [15,16]. Fig. 1 shows the shape of the D7-brane, which is the scalar field configuration $w(\rho)$, for $E/E_{\rm Sch} = 0.5, 0.8, 0.95, 1$ and the critical embedding. At the critical embedding which is a confinement/deconfinement phase boundary although the solution itself

[1] We solve the full nonlinear DBI action with the electric field while expanding the solution by a linear eigenbasis at the $E = 0$ vacuum, which is possible since the eigenbasis provides a complete set of the mutually orthogonal base modes for normalizable deformations. Alternatively, one can start with the meson effective Lagrangian which is written only in terms of the eigenmodes at the $E = 0$ vacuum. Then turning on the electric field makes the mesons get condensed, giving the same result.