



Thermal dissociation of heavy mesons and configurational entropy

Nelson R.F. Braga^a, Luiz F. Ferreira^a, Roldão da Rocha^{b,*}

^a Instituto de Física, Universidade Federal do Rio de Janeiro, Caixa Postal 68528, RJ 21941-972, Brazil

^b CMCC, Universidade Federal do ABC, UFABC, 09210-580, Santo André, Brazil

ARTICLE INFO

Article history:

Received 3 September 2018

Received in revised form 17 October 2018

Accepted 18 October 2018

Available online 22 October 2018

Editor: N. Lambert

ABSTRACT

We study the dissociation of heavy vector mesons in a thermal medium from the point of view of the configurational entropy. Bottomonium and charmonium vector mesons in a plasma are represented using a holographic AdS/QCD model. The quasinormal modes, describing the quasiparticle states, are determined for a representative range of temperature and their corresponding CE is then computed. The main result indicates that the CE is an important contrivance to investigate the stability of the heavy mesons against dissociation in the thermal medium.

© 2018 The Author(s). Published by Elsevier B.V. This is an open access article under the CC BY license (<http://creativecommons.org/licenses/by/4.0/>). Funded by SCOAP³.

1. Introduction

The configurational entropy (CE) is a quantity that implements the concept of informational entropy as an information logarithmic measure. The CE comprises the way how the information is compressed into any physical system [1,2]. The CE represents a measure of information that regards the spatial complexity of a localized system, describing the data that is necessary to depict the spatial shape of square-integrable scalar fields [3,4]. The modes in a system with lower CE have been shown to be more dominant, and hence more detectable and observable [5]. Further aspects of the CE were presented in Ref. [6].

Regarding quantum chromodynamics (QCD), describing how gluons and quarks strongly interact, the CE has been playing a prominent role together with holographic AdS/QCD models, to reveal important features regarding QCD phenomenology. For example, light flavor mesonic excitations [5] and scalar glueballs [7] were scrutinized using the CE, being further generalized in a dynamical holographic model setup [8]. Besides, the CE underlying bottomonium and charmonium at zero temperature confirmed the experimental predominance of lower S wave resonances and lower masses quarkonia states, in Ref. [9]. Still in the AdS/QCD paradigm, ρ and ϕ mesons were studied in Ref. [10], besides the color-glass condensate in the quark–gluon plasma (QGP) [11,12], both in the context of the CE. Moreover, solitonic solutions of the KdV equation, arising in a cold QGP were also studied with the instruments

of the CE, in Ref. [13]. The CE applications to heavy ion collisions (HICs) was scrutinized in Ref. [14].

The CE was shown to be, moreover, very useful as an apparatus to probe gravity and related phenomena. For example, the Hawking–Page phenomenon of phase transition leading thermal radiation into AdS–Schwarzschild black branes was investigated in Ref. [15] in the context of the CE. After the seminal applications of the CE to the study of stellar stability in Refs. [16,17], Bose–Einstein condensates of gravitons were also scrutinized, where the critical stellar density was derived as a global minimum of the CE, in Ref. [18]. The CE was further employed to investigate domain walls [19] and vortexes [20], and their applications in field theory [21,22]. Besides, spectral functions of bottomonium states were used in the context of the information entropy method [23].

HICs, produced in particle accelerators, lead to the formation of a QGP. This deconfined phase of hadronic matter survives only for very short time scales, being indirectly observed, analysing the collisions yields in the detectors. In particular, the fraction of heavy vector mesons detected can provide information about the plasma state. The suppression of such particles in the by products of a HIC, as compared to a proton–proton collision, serves as an indicator that a thermal medium was formed [24,25].

An interesting tool to describe heavy vector mesons is to use the AdS/QCD holographic models. Inspired on gauge string AdS/CFT duality, these models are particularly useful in the description of hadrons in a thermal medium. A holographic model for $b\bar{b}$ and $c\bar{c}$ vector mesons (bottomonium and charmonium respectively), involving two dimensionfull parameters, was proposed in Refs. [26,27]. Then, an improved holographic model that contains three energy parameters appeared in Refs. [28–30]. Two of these parameters represent the quark mass and the string ten-

* Corresponding author.

E-mail addresses: braga@if.ufrj.br (N.R.F. Braga), luizfaulhaber@if.ufrj.br (L.F. Ferreira), roldao.rocha@ufabc.edu.br (R. da Rocha).

sion, necessary in order to fit the mass spectra of quarkonium states. The third parameter represents an ultraviolet (UV) energy scale, needed to fit the spectra of decay constants. They are associated with non-hadronic decay, proportional to the transition matrix from a meson state at excitation level n to the hadronic vacuum, $\langle 0 | J_\mu(0) | n \rangle = \epsilon_\mu f_n m_n$. The UV energy parameter is related to the large mass change that occurs in such a process, when heavy vector mesons decay to light leptons.

At finite temperature, the mesons are described by the thermal spectral function, that composes the imaginary part of the retarded Green's function. The relevant part of the retarded Green's function is the 2-point function, that at zero temperature has a spectral decomposition in terms of masses m_n and decay constants f_n of the states,

$$\Pi(p^2) = \sum_{n=1}^{\infty} \frac{f_n^2}{(-p^2) - m_n^2 + i\epsilon}. \quad (1)$$

The imaginary part of Eq. (1) represents a sum of Dirac delta peaks, with coefficients that are proportional to the square of the decay constants, $f_n^2 \delta(-p^2 - m_n^2)$. At finite temperature, the peaks get smeared, as the temperature T of the medium increases. This connection between the spectral function and the decay constants is the reason why a consistent extension of a hadronic model to finite temperature must be consistent with decay constants. The results that emerge from Refs. [27–30] show that as the temperature increases, the heavy vector mesons dissociate in the medium. In this letter we want to see how this dissociation process is translated into the CE.

This paper is organized as follows: Sect. 2 is devoted to briefly introducing the motivation and the key points in the CE paradigm. The 2-point correlators and their power spectrum are shown to give rise to a correlation probability distribution, the modal fraction, that is the leading ingredient to compute the CE. In Sect. 3 the holographic model is introduced for heavy vector mesons, whose pure gauge action yields normalizable solutions, used for deriving the decay constants for charmonium and bottomonium states, whose quasinormal modes in the finite temperature case are derived and analyzed. Sect. 4 is then dedicated to the compute the CE underlying quarkonia, at finite temperature, as a function of the S-wave resonances excitation level and the temperature as well. Relevant results impact the probability of dissociation with respect to the temperature, distinguishing the charmonium from the bottomonium.

2. Configurational entropy and Shannon information entropy

Given a probability distribution, the CE, motivated by the Shannon's theory of information, is defined for discrete systems by $S[f] = -\sum_a p_a \log(p_a)$, standing for an upper limit of the information compression [2,4]. The CE, thus, comprises the informational bulk of all configurations in a physical system. If all the N system modes present a uniform probability distribution, then one can assume $p_a = 1/N$, yielding the CE to attain its maximum value $S = \log N$. The (differential) CE takes into account the continuous limit of this paradigm [6], whose equivalent features were also discussed in Ref. [5].

To introduce the CE of a physical system, spatial correlations must be taken into account, characterizing the correlation of the fluctuations of a localized, square integrable, $\rho(x)$ scalar field, that describes the system. Small perturbations can make the scalar field to fluctuate. The correlator, namely, the 2-point correlation function reads $G(r) = \int_{\mathbb{R}^d} \rho(\tilde{r}) \rho(r + \tilde{r}) d^d \tilde{r}$, where d denotes the spatial dimension. The CE measures the informational theory of

correlations, when constructing a probability distribution from the correlator. For it, the power spectrum $P(k)$ is taken into account as the correlator Fourier transform. The convolution theorem makes it possible to express the power spectrum as $P(k) \sim \|\int_{\mathbb{R}^d} \rho(r) e^{ik \cdot r} d^d r\|^2$, being normalizable. It corresponds to the uniform distribution, if $\rho(r)$ is a random distribution. Since the normalizable power spectrum satisfies the condition to be a continuous probability distribution associated to some continuous random variable, taking the Fourier transform $\rho(k) = \int_{\mathbb{R}^d} \rho(r) e^{-ik \cdot r} d^d x$, the correlation probability distribution, known as the modal fraction, is defined as [2]

$$\epsilon(k) = \frac{|\rho(k)|^2}{\int_{\mathbb{R}^d} |\rho(k)|^2 d^d k}. \quad (2)$$

The modal fraction measures the contribution of a given mode, k , to the power spectrum. As the power spectrum describes the fluctuation of the scalar field spatial profile, the modal fraction measures the contribution of a range of modes for the description of the shape of the scalar field. The amount of information to describe the spatial profile of ρ , in terms of the Fourier modes, is given by the CE,

$$S[\rho] = - \int_{\mathbb{R}^d} \epsilon_o(k) \log \epsilon_o(k) d^d k, \quad (3)$$

where $\epsilon_o(k) = \epsilon(k)/\epsilon_{\max}(k)$, being the denominator the maximum fraction. In the case where the modal fraction represents the uniform distribution, it maximizes the CE, corresponding to a random field for which the power spectrum is the uniform distribution.

3. Holographic model

For the zero temperature case, the model proposed in Refs. [29, 30] for charmonium and bottomonium, representing quarkonia states, is defined in a 5D anti-de Sitter (AdS) space–time

$$ds^2 = \frac{R^2}{z^2} (-dt^2 + dx^i dx_i + dz^2). \quad (4)$$

The heavy vector mesons are described by a vector field $V_m = (V_\mu, V_z)$ ($\mu = 0, 1, 2, 3$), assumed to be dual to the fermionic gauge theory current density $J^\mu = \bar{\psi} \gamma^\mu \psi$. The action for this field reads

$$I = \int d^4 x dz \sqrt{-g} e^{-\phi(z)} \left(-\frac{1}{4g_5^2} F^{mn} F_{mn} \right), \quad (5)$$

where $F_{mn} = \partial_m V_n - \partial_n V_m$ and $\phi(z)$ is a phenomenological background field with the form

$$\phi(z) = k^2 z^2 + Mz + \tanh \left(\frac{1}{Mz} - \frac{k}{\sqrt{\Gamma}} \right). \quad (6)$$

It is important to note that this is a phenomenological effective model for heavy vector mesons. The background that we use is not a solution of the Einstein's equations, since we included the scalar field of Eq. (6). This type of model is usually called a bottom up holographic model, as opposed to the so called top down models that use backgrounds that are solutions of Einstein's equations. The parameters play the following role: k represents the quark mass, Γ the string tension of the strong quark anti-quark interaction and M is a mass scale associated with non hadronic decay.

One chooses the gauge $V_z = 0$. Then, the equation of motion for the spatial transverse components of the field, denoted generically as V , takes in momentum space the following form,

Table 1

Holographic masses and decay constants for the charmonium S -wave resonances. Experimental values between parentheses for comparison.

Holographic (and experimental) results for charmonium		
State	Mass (MeV)	Decay constants (MeV)
1S	2943(3096.916 \pm 0.011)	399(416 \pm 5.3)
2S	3959(3686.109 \pm 0.012)	255(296.1 \pm 2.5)
3S	4757(4039 \pm 1)	198(187.1 \pm 7.6)
4S	5426(4421 \pm 4)	169(160.8 \pm 9.7)

Table 2

Holographic masses and decay constants for the bottomonium S -wave resonances. Experimental values between parentheses for comparison.

Holographic (and experimental) results for bottomonium		
State	Mass (MeV)	Decay constants (MeV)
1S	6905(9460.30 \pm 0.26)	719(715.0 \pm 2.4)
2S	8871(10023.26 \pm 0.32)	521(497.4 \pm 2.2)
3S	10442(10355.2 \pm 0.5)	427(430.1 \pm 1.9)
4S	11772(10579.4 \pm 1.2)	375(340.7 \pm 9.1)

$$\partial_z \left[e^{-B(z)} \partial_z V \right] - p^2 e^{-B(z)} V = 0, \quad (7)$$

where $B(z) = \log\left(\frac{z}{R}\right) + \phi(z)$.

The meson states at zero temperature are represented by the normalizable solutions, $V(p, z) = \Psi_n(z)$, that satisfy the boundary conditions $\Psi_n(0) = 0$, for $p^2 = -m_n^2$, where m_n denotes the masses of the corresponding meson states. These eigenfunctions are normalized as $\int_0^\infty e^{-B(z)} \Psi_n(z) \Psi_m(z) dz = \delta_{mn}$, and the decay constants are calculated holographically as

$$f_n = \frac{1}{g_5 m_n} \lim_{z \rightarrow 0} \left(e^{-B(z)} \partial_z \Psi_n(z) \right). \quad (8)$$

The values of the parameters in the background (6), obtained in Ref. [30] in order to describe charmonium and bottomonium, are respectively:

$$k_c = 1.2 \text{ GeV}, \quad \sqrt{\Gamma_c} = 0.55 \text{ GeV}, \quad M_c = 2.2 \text{ GeV}, \quad (9)$$

$$k_b = 2.45 \text{ GeV}, \quad \sqrt{\Gamma_b} = 1.55 \text{ GeV}, \quad M_b = 6.2 \text{ GeV}. \quad (10)$$

Tables 1 and 2 show the results for charmonium and bottomonium, respectively. For comparison, the experimental data from Ref. [31] is show between parentheses. It is worth to mention that the higher the radial excitation level, the lower the decay constants are.

For the finite temperature case, we consider the same action of Eq. (5), now in an AdS black hole space-time

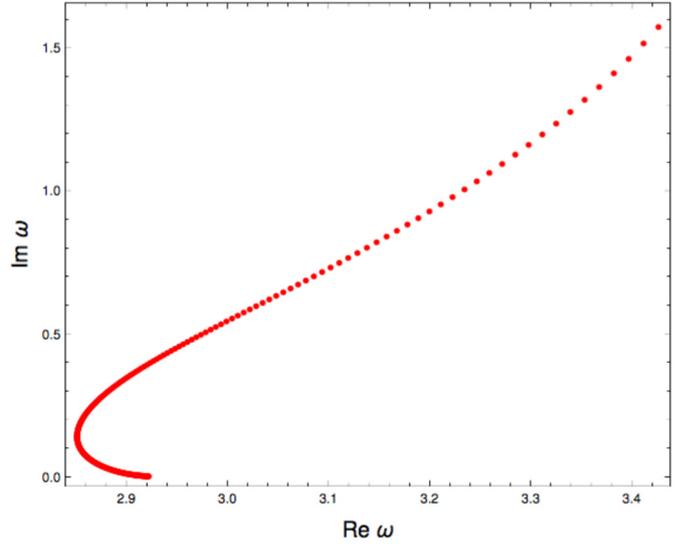
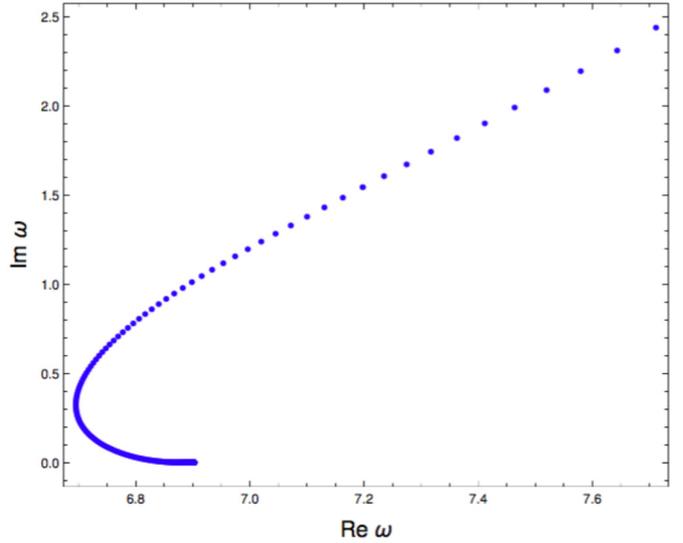
$$ds^2 = \frac{R^2}{z^2} \left(-f(z) dt^2 + \frac{dz^2}{f(z)} + dx^i dx_i \right), \quad (11)$$

where

$$f(z) = 1 - \frac{z^4}{z_h^4}. \quad (12)$$

The horizon position z_h is determined from the condition $f(z_h) = 0$ and the black hole temperature comes from the requirement of absence of a conical singularity at the horizon, in the Euclidean version of the metric. This condition implies a relation between the periodic time interval $0 \leq t \leq \beta = 1/T$ and the horizon position, given by

$$T = \frac{|f'(z)|_{(z=z_h)}}{4\pi} = \frac{1}{\pi z_h}. \quad (13)$$

**Fig. 1.** Quasinormal modes for charmonium.**Fig. 2.** Quasinormal modes for bottomonium.

In contrast to the zero temperature case, where the particles are described by normalized solutions of the field, at finite temperature the quasiparticles are described by solutions known as quasinormal modes. Previous studies of quasinormal modes in the context of gauge/gravity duality can be found, for example, in Refs. [33,32,34–42]. Quasinormal modes are obtained solving the equation that comes from the metric (11):

$$\partial_z (f(z) e^{-B(z)} \partial_z V) + \frac{\omega^2}{f(z)} e^{-B(z)} V = 0, \quad (14)$$

with the following boundary conditions:

$$v_n(0) = 0, \quad (15)$$

$$\lim_{z \rightarrow z_h} v_n = \left(1 - \frac{z}{z_h} \right)^{-i\omega/(4\pi T)} \left[1 + a_1 \left(1 - \frac{z}{z_h} \right) + \dots \right].$$

The second condition corresponds to selecting just infalling boundary conditions at the horizon position. The coefficients that appear in the infalling condition can be determined inserting this condition into Eq. (14). Similarly to normal modes, the quasinormal

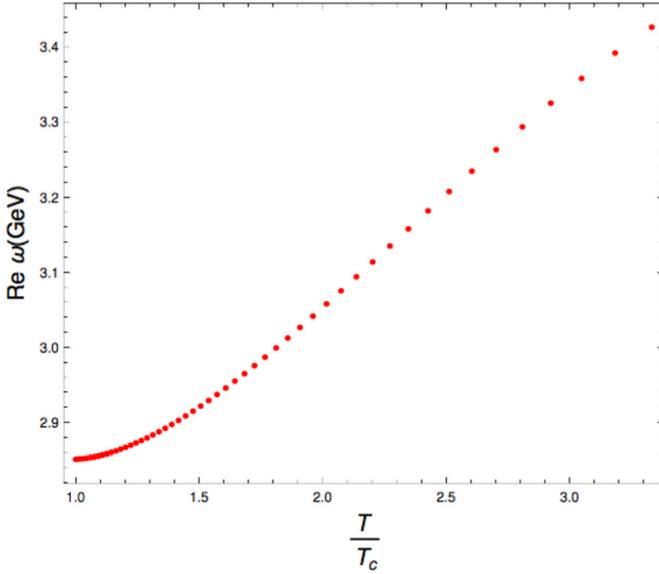


Fig. 3. Real part of the frequency for charmonium quasinormal modes.

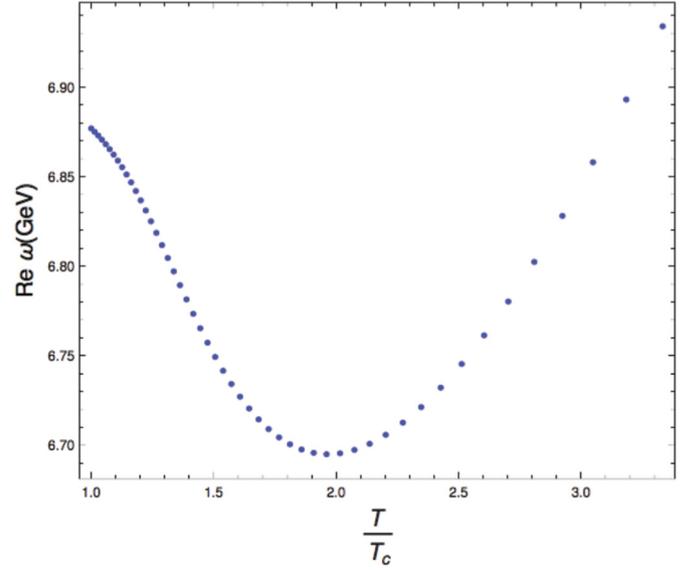


Fig. 5. Real part of the frequency for bottomonium quasinormal modes.

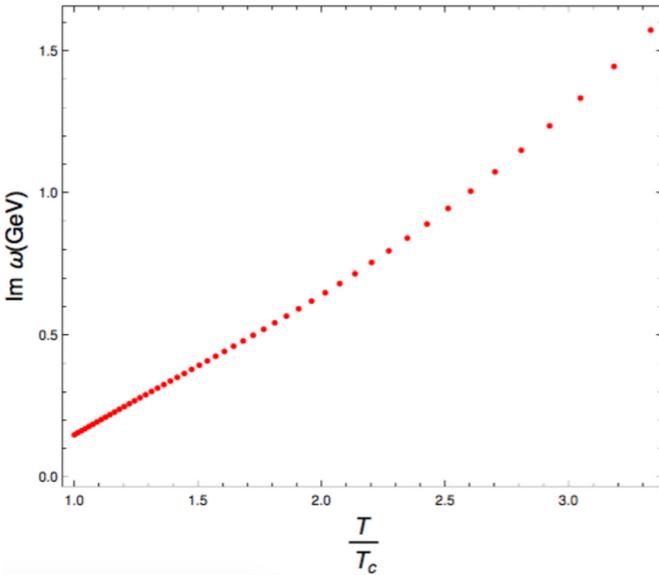


Fig. 4. Imaginary part of the frequency for charmonium quasinormal modes.

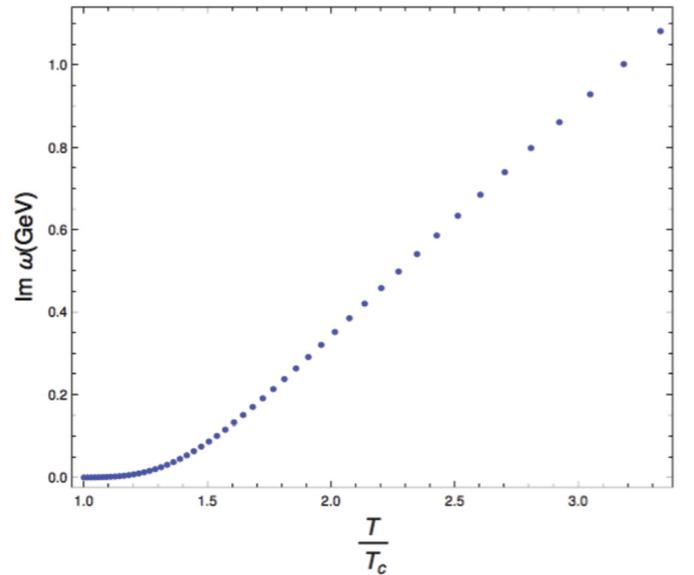


Fig. 6. Imaginary part of the frequency for bottomonium quasinormal modes.

modes exist only for a discrete set of frequencies ω_n . However, in the case of quasinormal modes, these frequencies are complex.

The method used to find the quasinormal modes in this work was the shooting method [43–45], whose basic idea consists of specifying two boundary conditions at the horizon, and then adjust the free parameter given by the frequency ω , to find the desired solution at the boundary. One of these two boundary conditions is the infalling boundary condition, presented in Eq. (15). The second boundary condition is the derivative of the infalling boundary condition. It is necessary to ensure that the solution found by solving Eq. (14), using the boundary conditions at the horizon, is a quasinormal mode varying the frequency ω until the field vanishes at the boundary.

The quasinormal modes of charmonium and bottomonium are determined as function of the temperature. We show in Figs. 1 and 2 the results for the imaginary part of the frequencies as a function of the real part for the first mode $n = 0$. In Figs. 3–6, the real and imaginary parts of the quasinormal modes are plotted for both

the charmonium and bottomonium respectively, as functions of the temperature. One notes that in the zero temperature limit, the real part of the quasinormal modes coincides with the corresponding quarkonium mass.

4. Configurational entropy of quarkonia at finite temperature

We now apply the CE paradigm to the analysis of heavy mesons. The action integral of the holographic model that we use is not consistent with the Einstein's equations, since we include the phenomenological scalar background of Eq. (6). However, we will assume that the energy momentum tensor of the model is obtained from the action in the same way as in general relativity. For an action integral of the form $\int d^4x dz \sqrt{-g} \mathcal{L}$ the energy momentum tensor reads

Table 3
Configurational entropy (CE) of bottomonium with respect to the S-wave resonances excitation level.

Bottomonium		
S-wave resonance	Masses (GeV)	CE
1S	6.9051	2.1786
2S	8.8710	2.4957
3S	10.442	2.6899
4S	11.772	2.8188
5S	12.943	2.9185
6S	14.001	3.0008
7S	14.972	3.0712

$$T_{mn}(z) = \frac{2}{\sqrt{-g}} \left[\frac{\partial(\sqrt{-g}\mathcal{L})}{\partial g^{mn}} - \frac{\partial}{\partial x^p} \frac{\partial(\sqrt{-g}\mathcal{L})}{\partial \left(\frac{\partial g^{mn}}{\partial x^p}\right)} \right] \quad (16)$$

Note that in the case of the action (5), the second term of (16) is absent, since the action does depend on derivatives of the metric. The energy density $\rho(z) = T_{00}(z)$ for the case of a vector field reads

$$\rho(z) = \frac{e^{-\phi(z)}}{g_5^2} \left[g_{00} \left(\frac{1}{4} g^{mp} g^{nq} F_{mn} F_{pq} \right) - g^{mn} F_{0n} F_{0m} \right]. \quad (17)$$

Considering the metric (11) and taking a plane wave solution in the x^μ directions, in the meson rest frame $V_\mu = \epsilon_\mu v(p, z) e^{-i\omega t}$, with $\epsilon_\mu = (0, 1, 0, 0)$, the energy density takes the form

$$\rho(z) = \frac{e^{-\phi(z)}}{2g_5^2 R^2} \left[- \left(1 - \frac{z^4}{z_h^4} \right)^2 (\partial_z v)^2 + |\omega|^2 v^2 \right] \quad (18)$$

Note that for complex solutions, the field strength term of the Lagrangian density takes the form $F_{mn}^* F^{mn}$. In Figs. 1–6, we present the real and imaginary parts of the frequency as a function of the temperature. The solutions for $v(p, z)$ at finite temperature are the quasinormal modes $v_n(z)$, described in the previous section. In the zero temperature limit, these solutions become the eigenfunctions of Eq. (7). Taking a Fourier transform in coordinate z and splitting $\rho(k) = C(k) + iS(k)$, where

$$C(k) = \int_0^\infty \rho(z) \cos(kz) dz, \quad (19)$$

$$S(k) = \int_0^\infty \rho(z) \sin(kz) dz, \quad (20)$$

one finds the modal fraction (2)

$$\epsilon(k) = \frac{S^2(k) + C^2(k)}{\int_0^\infty [S^2(\dot{k}) + C^2(\dot{k})] d\dot{k}}. \quad (21)$$

Finally, the CE (3) reads

$$S[\rho] = - \int_0^\infty \epsilon_o(k) \log \epsilon_o(k) dk. \quad (22)$$

Using the appropriate values of the parameters for the quarkonia (9)–(10) and employing Eqs. (3) and (18), we compute the CE for the quasi-states. First one evaluates the CE for the states at zero temperature using the eigenfunctions of Eq. (7). The results are listed in Tables 3 and 4. Then, the results for the 1S quarkonia quasi-states, at finite temperature, are presented in Fig. 7 as a

Table 4
Configurational entropy (CE) with respect to the S-wave resonances excitation level.

Charmonium		
S-wave resonance	Masses (GeV)	CE
1S	2.9425	1.4105
2S	3.9586	1.7413
3S	4.7573	1.9457
4S	5.4261	2.0793
5S	6.0114	2.1791
6S	7.0201	2.3130
7S	7.4675	2.4526

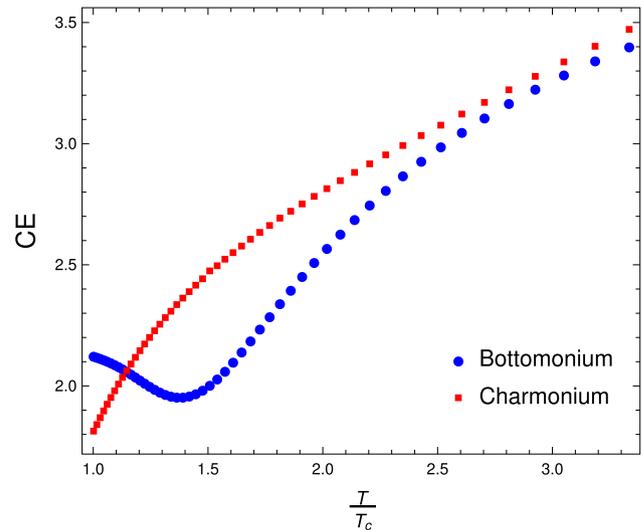


Fig. 7. Configurational entropy (CE) for bottomonium and charmonium 1S quasi-state as a function of the ratio T/T_c .

function of T/T_c , where T_c is the critical deconfinement temperature at which the deconfined medium is formed. Using lattice QCD studies, Refs. [46–49] show that heavy mesons only melt at higher temperatures than T_c .

From Fig. 7, one notes that, for the charmonium states, the CE monotonically increases with the temperature. An increment in the CE is expected to represent a decrease in the stability of the system. Charmonium quasi-states were studied in Ref. [30] through the determination of the spectral function. The 1S quasi-state appears as a peak that gets continuously smeared as the temperature increases. The higher the temperature, the lower and broader is the peak, corresponding to a continuous increment in the degree of dissociation in the medium. The quasi-normal modes obtained in Sect. 3 show an equivalent behaviour. The imaginary part of the quasi normal mode frequency increases with the temperature. This imaginary component is related to the width of the peak. So, the continuous increment in this quantity represents the increase of the dissociation degree. Charmonium states survive the deconfinement transition, in contrast to light hadrons that completely dissociate. But, as seen from the spectral function and from the quasinormal modes, they undergo a partial dissociation, with a degree that increases with the temperature. The result obtained for the CE is consistent with this thermal behaviour. In fact, the lower the CE, the more informationally stable the state is, corresponding to a bigger dominance of the respective state in the system. The probability of dissociation in the medium continuously increases with the temperature.

For charmonium states, already in Fig. 7 one can realize a strong correlation between the CE and the temperature T/T_c ratio. The higher the temperature, the higher the charmonium CE is. Besides, entropic Regge-like trajectories can be implemented

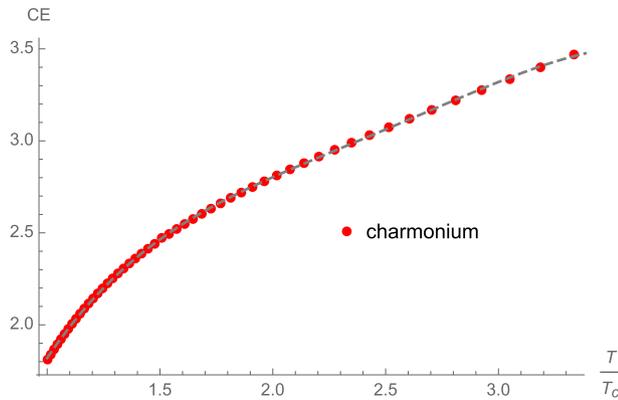


Fig. 8. CE with respect to the temperature T/T_c ratio, for the charmonium, and its quartic interpolation.

in the context of AdS/QCD, that associates the charmonium CE with the temperature T/T_c ratio. In fact, interpolating the charmonium CE we find $CE\left(\frac{T}{T_c}\right) = -0.1008\frac{T^4}{T_c^4} + 1.0103\frac{T^3}{T_c^3} - 3.7726\frac{T^2}{T_c^2} + 6.7419\frac{T}{T_c} - 2.0590$, at the analyzed temperature T/T_c ratio range, within 0.08%. It indicates a scaling law between the CE and the temperature T/T_c ratio, considering all the range $1 \lesssim T/T_c \lesssim 3.5$.

For the bottomonium 1S state the result is different. Indeed, the CE decreases up to temperatures of the order of $\sim 1.3T_c$ and then increases for higher temperatures. Comparing with results from the analysis of the spectral function in Ref. [30], one should expect an increasing dissociation degree as the temperature increases. Such a behaviour should be translated into a monotonic increase in the CE, although the bottomonium state survives at temperatures much higher than the charmonium. It is important to note, when comparing charmonium and bottomonium results, that the holographic model provides results with different accuracy degrees for the two cases. From Tables 1 and 2, one sees that for the charmonium 1S state the result for the mass differs just 5% from the experimental value. In contrast, for the bottomonium 1S state the mass of the model differs 27% from the experimental value. It is a very difficult task to find a simple holographic model that fits masses and decay constants for bottomonium. It is possible that this problem with the fit of the value of the mass is leading to the failure of a consistent description of bottomonium dissociation, in terms of configurational entropy, for low temperatures $T \lesssim 1.3T_c$.

The model represents a thermal medium, the plasma, that is present only for temperatures above the critical one T_c . We considered, when plotting Fig. 7, that the critical temperature is $T_c = 191$ MeV [50], corresponding to the transition temperature, when the model is used for light flavor mesons.

A scaling law, between the CE and the temperature T/T_c ratio, also exists for the bottomonium state for temperatures higher than $\sim 1.35T_c$. The quartic polynomial $CE\left(\frac{T}{T_c}\right) = 0.2947\frac{T^4}{T_c^4} - 2.8530\frac{T^3}{T_c^3} + 9.8565\frac{T^2}{T_c^2} - 13.5219\frac{T}{T_c} + 8.2633$, at the range $1.35 \lesssim T/T_c \lesssim 3.5$, fits the graphic for the bottomonium CE, within 0.29%. This is plotted in Fig. 9.

5. Conclusions

The dissociation of heavy vector mesons, in the framework of holographic AdS/QCD, was studied in the context of the configurational entropy setup. Charmonium and bottomonium states were analyzed from the point of view of their quasinormal modes, at finite temperature. The quarkonia quasinormal modes are plotted in Figs. 1–6. The CE underlying quarkonia was analyzed both as a

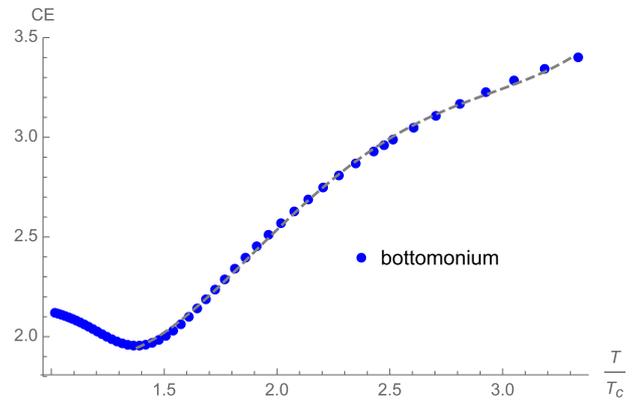


Fig. 9. CE with respect to the temperature T/T_c ratio, for the bottomonium, and its quartic interpolation.

function of the temperature (Figs. 7–9) and of the radial excitation level of S-wave resonances (Tables 3 and 4).

Important results make the quarkonia probability of dissociation to be distinct, with respect to the temperature. In fact, Figs. 7–9 show the CE of bottomonium and charmonium as a function of the temperature T/T_c ratio. The charmonium states were shown to present their underlying CE that increases monotonically with the temperature. In this case, the probability of dissociation in the medium increases with the temperature. For the bottomonium, Figs. 7 and 9 show a quite different result. First, the CE computed for the bottomonium has a global minimum, corresponding to temperatures $T \sim 1.3T_c$. Solely for higher temperatures, the CE increases monotonically. One feasible explanation for this phenomenon is that the bottomonium energy parameters are higher, compared to the charmonium ones. The AdS/QCD holographic model at finite temperature is assumed to represent a plasma when $T \gtrsim T_c$, at which the deconfined medium is formed. The results for bottomonium indicate that only for $T \gtrsim 1.35T_c$ the model presents an appropriate description of the thermal deconfined medium. It was possible to find for charmonium and bottomonium states, an interpolation of the CE as a function of the temperature T/T_c ratio. Figs. 8 and 9 respectively show, for the charmonium and the bottomonium, scaling laws relating the CE to the temperature T/T_c ratio. Our results indicate that the CE is a relevant tool to investigate stability of the heavy mesons against dissociation in the thermal medium.

Acknowledgements

N.B. is partially supported by CNPq and L. F. Ferreira is supported CAPES; RdR is grateful to Dr. A. Gonçalves for fruitful discussions and to FAPESP (Grant No. 2017/18897-8) and to the National Council for Scientific and Technological Development – CNPq (Grant No. 303293/2015-2), for partial financial support.

References

- [1] M. Gleiser, N. Stamatopoulos, *Phys. Lett. B* 713 (2012) 304, arXiv:1111.5597 [hep-th].
- [2] M. Gleiser, N. Stamatopoulos, *Phys. Rev. D* 86 (2012) 045004, arXiv:1205.3061 [hep-th].
- [3] M. Gleiser, N. Graham, *Phys. Rev. D* 89 (2014) 083502, arXiv:1401.6225 [astro-ph.CO].
- [4] M. Gleiser, D. Sowiński, *Phys. Lett. B* 747 (2015) 125, arXiv:1501.06800 [cond-mat.stat-mech].
- [5] A.E. Bernardini, R. da Rocha, *Phys. Lett. B* 762 (2016) 107, arXiv:1605.00294 [hep-th].
- [6] M. Gleiser, M. Stephens, D. Sowiński, *Phys. Rev. D* 97 (9) (2018) 096007, arXiv:1803.08550 [hep-th].

- [7] A.E. Bernardini, N.R.F. Braga, R. da Rocha, Phys. Lett. B 765 (2017) 81, arXiv:1609.01258v1 [hep-th].
- [8] N. Barbosa-Cendejas, R. Cartas-Fuentevilla, A. Herrera-Aguilar, R.R. Mora-Luna, R. da Rocha, Phys. Lett. B 782 (2018) 607, arXiv:1805.04485 [hep-th].
- [9] N.R.F. Braga, R. da Rocha, Phys. Lett. B 776 (2018) 78, arXiv:1710.07383 [hep-th].
- [10] G. Karapetyan, Phys. Lett. B 781 (2018) 201, arXiv:1802.09105 [nucl-th].
- [11] G. Karapetyan, Europhys. Lett. 117 (2017) 18001, arXiv:1612.09564 [hep-ph].
- [12] G. Karapetyan, Europhys. Lett. 118 (2017) 38001, arXiv:1705.10617 [hep-ph].
- [13] A. Goncalves, R. da Rocha, Phys. Lett. B 774 (2017) 98, arXiv:1706.01482 [hep-ph].
- [14] C.W. Ma, Y.G. Ma, Prog. Part. Nucl. Phys. 99 (2018) 120, arXiv:1801.02192 [nucl-th].
- [15] N.R.F. Braga, R. da Rocha, Phys. Lett. B 767 (2017) 386, arXiv:1612.03289 [hep-th].
- [16] M. Gleiser, D. Sowiński, Phys. Lett. B 727 (2013) 272, arXiv:1307.0530 [hep-th].
- [17] M. Gleiser, N. Jiang, Phys. Rev. D 92 (2015) 044046, arXiv:1506.05722 [gr-qc].
- [18] R. Casadio, R. da Rocha, Phys. Lett. B 763 (2016) 434.
- [19] R.A.C. Correa, R. da Rocha, Eur. Phys. J. C 75 (2015) 522, arXiv:1502.02283 [hep-th].
- [20] R.A.C. Correa, D.M. Dantas, C.A.S. Almeida, R. da Rocha, Phys. Lett. B 755 (2016) 358, arXiv:1601.00076 [hep-th].
- [21] A. Alves, A.G. Dias, R. da Silva, Physica 420 (2015) 1, arXiv:1408.0827 [hep-ph].
- [22] A. Alves, A.G. Dias, R. da Silva, Braz. J. Phys. 47 (2017) 426.
- [23] G. Aarts, C. Allton, T. Harris, S. Kim, M.P. Lombardo, S.M. Ryan, J.I. Skullerud, J. High Energy Phys. 1407 (2014) 097, arXiv:1402.6210 [hep-lat].
- [24] T. Matsui, H. Satz, Phys. Lett. B 178 (1986) 416.
- [25] H. Satz, J. Phys. G 32 (2006) R25, arXiv:hep-ph/0512217.
- [26] N.R.F. Braga, M.A. Martin Contreras, S. Diles, Phys. Lett. B 763 (2016) 203, arXiv:1507.04708 [hep-th].
- [27] N.R.F. Braga, M.A. Martin Contreras, S. Diles, Eur. Phys. J. C 76 (11) (2016) 598, arXiv:1604.08296 [hep-ph].
- [28] N.R.F. Braga, L.F. Ferreira, Phys. Lett. B 773 (2017) 313, arXiv:1704.05038 [hep-ph].
- [29] N.R.F. Braga, L.F. Ferreira, A. Vega, Phys. Lett. B 774 (2017) 476, arXiv:1709.05326 [hep-ph].
- [30] N.R.F. Braga, L.F. Ferreira, Phys. Lett. B 783 (2018) 186, arXiv:1802.02084 [hep-ph].
- [31] K.A. Olive, et al., Particle Data Group Collaboration, Chin. Phys. C 38 (2014) 090001.
- [32] G.T. Horowitz, V.E. Hubeny, Phys. Rev. D 62 (2000) 024027, arXiv:hep-th/9909056.
- [33] A. Nunez, A.O. Starinets, Phys. Rev. D 67 (2003) 124013, arXiv:hep-th/0302026.
- [34] P.K. Kovtun, A.O. Starinets, Phys. Rev. D 72 (2005) 086009, arXiv:hep-th/0506184.
- [35] A.S. Miranda, V.T. Zanchin, Phys. Rev. D 73 (2006) 064034, arXiv:gr-qc/0510066.
- [36] C. Hoyos-Badajoz, K. Landsteiner, S. Montero, J. High Energy Phys. 0704 (2007) 031, arXiv:hep-th/0612169.
- [37] E. Berti, V. Cardoso, A.O. Starinets, Class. Quantum Gravity 26 (2009) 163001, arXiv:0905.2975 [gr-qc].
- [38] J. Morgan, V. Cardoso, A.S. Miranda, C. Molina, V.T. Zanchin, Phys. Rev. D 80 (2009) 024024, arXiv:0906.0064 [hep-th].
- [39] A.S. Miranda, C.A. Ballon Bayona, H. Boschi-Filho, N.R.F. Braga, J. High Energy Phys. 0911 (2009) 119, arXiv:0909.1790 [hep-th].
- [40] R.A. Konoplya, A. Zhidenko, Rev. Mod. Phys. 83 (2011) 793, arXiv:1102.4014 [gr-qc].
- [41] L.A.H. Mamani, A.S. Miranda, H. Boschi-Filho, N.R.F. Braga, J. High Energy Phys. 1403 (2014) 058, arXiv:1312.3815 [hep-th].
- [42] L.A.H. Mamani, J. Morgan, A.S. Miranda, V.T. Zanchin, Phys. Rev. D 98 (2) (2018) 026006, arXiv:1804.01544 [gr-qc].
- [43] M. Kaminski, Fortschr. Phys. 57 (2009) 3, arXiv:0808.1114 [hep-th].
- [44] M. Kaminski, K. Landsteiner, F. Pena-Benitez, J. Erdmenger, C. Greubel, P. Kerner, J. High Energy Phys. 1003 (2010) 117, arXiv:0911.3544 [hep-th].
- [45] S. Janiszewski, M. Kaminski, Phys. Rev. D 93 (2) (2016) 025006, arXiv:1508.06993 [hep-th].
- [46] G. Aarts, C. Allton, M.B. Oktay, M. Peardon, J.I. Skullerud, Phys. Rev. D 76 (2007) 094513, arXiv:0705.2198 [hep-lat].
- [47] M. Asakawa, T. Hatsuda, Phys. Rev. Lett. 92 (2004) 012001, arXiv:hep-lat/0308034.
- [48] M. Fujita, K. Fukushima, T. Misumi, M. Murata, Phys. Rev. D 80 (2009) 035001, arXiv:0903.2316 [hep-ph].
- [49] M. Fujita, T. Kikuchi, K. Fukushima, T. Misumi, M. Murata, Phys. Rev. D 81 (2010) 065024, arXiv:0911.2298 [hep-ph].
- [50] C.P. Herzog, Phys. Rev. Lett. 98 (2007) 091601, arXiv:hep-th/0608151.