Meson widths from string worldsheet instantons

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ABSTRACT

We show that open strings living on a D-brane which lies outside an AdS black hole can tunnel into the black hole through worldsheet instantons. These instantons have a simple interpretation in terms of thermal quarks in the dual Yang–Mills (YM) theory. As an application we calculate the width of a meson in a strongly coupled quark–gluon plasma which is described holographically as a massless mode on a D7 brane in $AdS_5 \times S_5$. While the width of the meson is zero to all orders in the $1/\sqrt{\lambda}$ expansion with λ the 't Hooft coupling, it receives non-perturbative contributions in $1/\sqrt{\lambda}$ from worldsheet instantons. We find that the width increases quadratically with momentum at large momentum and comment on potential phenomenological implications of this enhancement for heavy ion collisions. We also comment on how this non-perturbative effect has important consequences for the phase structure of the YM theory obtained in the classical gravity limit.

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A heavy quarkonium bound state, like J/ψ or Υ, which finds itself in the quark–gluon plasma (QGP), becomes increasingly unstable and eventually dissociates at sufficiently high temperatures. On the one hand, this can be attributed to the weakening attraction between a heavy quark and antiquark in the bound state due to color screening of the medium [1]. On the other hand, the bound state can be broken up from collisions with the deconfined quarks and gluons in the medium [2]. Given that the QGP at RHIC and LHC is likely strongly coupled, understanding such medium effects on the propagation and dissociation of heavy quarkonia presents nontrivial theoretical challenges which must be confronted in order to interpret experimental data on quarkonium suppression.

Interesting insights have recently been made into this problem from strongly coupled model gauge theories using the AdS/CFT correspondence. AdS/CFT-based methods are powerful at attacking questions of dynamical origin, such as how the motion of quarkonia relative to the medium affects their various properties. The simplest example of the correspondence is provided by the duality between $\mathcal{N} = 4$ $SU(N_c)$ super-Yang–Mills (SYM) theory and string theory in $AdS_5 \times S_5$ [3]. Based on a calculation of the potential between a heavy external quark and antiquark pair moving in the strongly coupled hot $\mathcal{N} = 4$ plasma, it has been argued in [4] (see also [5]) that the dissociation temperature T_{diss} of a heavy quarkonium *decreases* with their velocity v relative to the medium as $T_{\text{diss}}(v) \approx (1 - v^2)^{1/4} T_{\text{diss}}(v = 0)$. Such a velocity scaling, which can be heuristically understood as due to increased screening from the boosted medium, could provide a significant additional source

of quarkonium suppression at nonzero transverse momentum in heavy ion collisions [4].

Rather than drawing inferences from the heavy quark potential, it is also possible to directly study the propagation of mesons in a strongly coupled plasma. While $\mathcal{N} = 4$ SYM theory itself does not contain dynamical mesons, one can obtain a closely related theory which does contain mesons by adding to it $N_f \ll N_c$ fundamental "quarks", which corresponds to adding some D7-branes to $AdS_5 \times S_5$ in the gravity picture [6]. It was found in [7,8] that meson dispersion relations are dramatically modified by the plasma and in particular, there exists a limiting velocity $v_c(T) < 1$, which decreases with increasing temperature. The existence of a subluminal limiting velocity is consistent with the velocity-enhanced screening obtained from the heavy quark potential, as when $v > v_c(T)$ the quark and antiquark are completely screened and no bound states are possible.

For a more complete understanding of the dissociation of mesons one needs to study their widths. We will be particularly interested in the momentum dependence of the widths. This has not been possible within the classical gravity approximation developed so far. In this approximation, the mesons are stable (i.e. they have zero width) for T smaller than a dissociation temperature T_{diss}, but completely disappear for $T > T_{\text{diss}}$ [7,9,10]. The approximation also has another important drawback: the densities of quarks and antiquarks are identically zero for a range of temperatures and chemical potentials [11] even though they should obey the standard thermal distribution.

In this Letter, we discuss a novel tunneling effect on the string worldsheet which gives rise to nonzero quark densities and meson widths for $T < T_{\text{diss}}$. This enables us to calculate explicitly the momentum dependence of the width of a meson in a strongly coupled

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QGP. We find that the width increases quadratically with momentum at large momentum.

At finite temperature, $\mathcal{N} = 4$ SYM theory can be described by a string theory in the spacetime of a black hole in $AdS_5 \times S_5$, whose metric can be written as

$$ds^2 = -f\,dt^2 + \frac{1}{f}\,dr^2 + \frac{r^2}{R^2}\,d\vec{x}^2 + R^2\,d\Omega_5^2, \qquad (1)$$

where $f = (r^2/R^2)(1 - r_0^4/r^4)$, $\vec{x} = (x_1, x_2, x_3)$. $d\Omega_5^2$ is the metric on a unit five-sphere S_5 which can be written as

$$d\Omega_5^2 = d\theta^2 + \sin^2\theta\,d\Omega_3^2 + \cos^2\theta\,d\phi^2, \quad \theta \in \left[0, \frac{\pi}{2}\right], \qquad (2)$$

with $d\Omega_3^2$ the metric for a three-sphere S_3. The string coupling g_s is related to the Yang–Mills coupling g_{YM} by $g_s = 4\pi g_{\text{YM}}^2$ and the curvature radius R is related to the 't Hooft coupling $\lambda = g_{\text{YM}}^2 N_c$ by $\frac{R^2}{\alpha'} = \sqrt{\lambda}$. The perturbative g_s and α' expansions in the bulk string theory are related to the $1/N_c$ and $\frac{1}{\sqrt{\lambda}}$ expansions in the Yang–Mills theory respectively. The temperature T of the YM theory is given by the Hawking temperature of the black hole, $T = \frac{r_0}{\pi R^2}$. Adding N_f fundamental "quarks" can be described in the dual string theory by adding N_f D7-branes in (1) [6]. A fundamental "quark" in the YM theory can be described by an open string with one end on the D7-branes and the other end on the black hole. Open strings with both ends on the D7-branes can be considered as "bound states" of a quark and antiquark, thus describing meson-type excitations in the YM theory.

We now briefly outline the standard procedure for obtaining the meson spectrum [9]. We will take $N_f = 1$, $N_c \to \infty$, and λ large but finite throughout the Letter. The D7-brane can be chosen to lie along the directions $\xi^\alpha = (t, \vec{x}, \Omega_3, \theta)$ and using the symmetries of the problem the embedding in the two remaining transverse directions can be taken as $\phi(\xi^\alpha) = 0$ and $r(\xi^\alpha) = r(\theta)$. At the lowest order in the α' expansion, $r(\theta)$ can be determined by extremizing the Dirac–Born–Infeld (DBI) action of the D7-brane with the boundary condition $r(\theta)\cos\theta|_{\theta \to \pi/2} \to L$, where L is related to the mass m_q of a quark in the Yang–Mills theory as $m_q = \frac{L}{2\pi\alpha'}$. For T smaller than some T_{diss}, $r(\theta)$ has the form shown in Fig. 1. The brane is closest to the black hole at $\theta = 0$, where there lies a 4-dimensional subspace spanned by $(t, \vec{x})$ since here the S_3 in (2) shrinks to a point. Denoting $r_m \equiv r(\theta = 0) > r_0$, the shortest open string connecting the D7-brane to the horizon has a mass in the YM theory

$$m_q^{(T)} = \frac{r_m - r_0}{2\pi\alpha'} = \sqrt{\lambda}\,T\,\frac{\Lambda_m - 1}{2}, \quad \Lambda_m = \frac{r_m}{r_0}. \qquad (3)$$

Note that Λ_m is a dimensionless number of $O(\lambda^0)$ determined by the ratio L/r_0, and $m_q^{(T)}$ can be interpreted as the effective mass of a quark at temperature T.

The mesons corresponding to massless fluctuations on the D7-brane can be found by solving the linearized equations resulting from expanding the DBI action around the embedding. For example, the quadratic action for the fluctuation $\chi^\phi(\xi^\alpha)$ of the location of D7-brane in the ϕ direction can be written as

$$S_{\text{DBI}}[\chi^\phi] = -\frac{\mu_7}{2} \int d^8\xi\,\sqrt{-g}\,G_{\phi\phi}\,g^{\alpha\beta}\,\partial_\alpha\chi^\phi\,\partial_\beta\chi^\phi, \qquad (4)$$

where $\mu_7 = \frac{1}{(2\pi)^7 g_s \alpha'^4}$ is the tension of the D7-brane, $G_{\phi\phi} = R^2\cos^2\theta$, and $g_{\alpha\beta}$ denotes the induced metric on the D7-brane. Writing $\chi^\phi = e^{-i\omega t + i\vec{k}\cdot\vec{x}}Y_l(\Omega_3)\psi(\theta)$, the equation of motion for ψ can be written as

$$\hat{H}(\vec{k}, l)\psi(\theta) = \omega^2\psi(\theta), \qquad (5)$$

where $\hat{H}(\vec{k}, l)$ is a second order differential operator in θ and Y_l are spherical harmonics on the S_3. For a given $\vec{k}, l$, $\hat{H}(\vec{k}, l)$ has
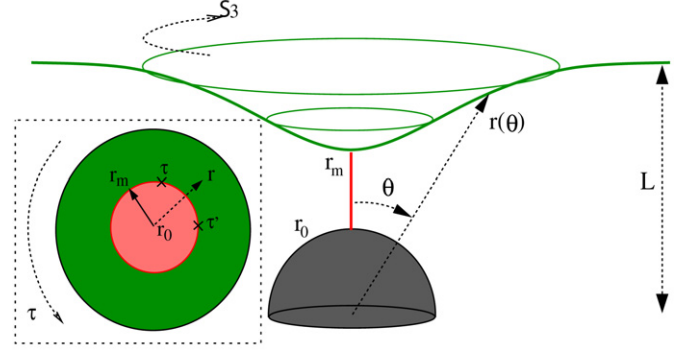

Fig. 1. An embedding of the D7 brane (green/upper curve) in the $AdS_5 \times S_5$ black hole geometry for $T < T_{\text{diss}}$ which lies entirely outside the black hole. Inset: the Euclidean r–τ plane at $\theta = 0$ showing a world-sheet instanton (red/inner disk) connecting the tip of D7 brane $r = r_m$ to the horizon at the center of the disk $r = r_0$. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this Letter.)

only discrete eigenvalues ω_n^2 labeled by an integer n, giving rise to dispersion relations $\omega = \omega_n(\vec{k}, l)$, all of which have zero width. In particular, the meson masses are of order $M = \frac{2\sqrt{2}L}{R^2} = \frac{4\sqrt{2\pi}m_q}{\sqrt{\lambda}}$. Since M is parametrically smaller than m_q in $\sqrt{\lambda}$, the mesons have a large binding energy, given by $2m_q^{(T)}$. There exists a temperature $T_{\text{diss}} = 0.122M$, beyond which the D7 brane falls into the black hole and mesons cease to exist as well-defined quasi-particles [7,10].

We stress that the zero-width conclusion only depends on the topology of the embedding in Fig. 1. Since mesons can only dissociate by falling into the black hole, when the D7-brane lies above the black hole horizon the mesons are necessarily stable. Given that the brane embedding and the background geometry are smooth, including higher order perturbative corrections in α' should not change the topology of the brane embedding if the distance between the brane and the horizon is parametrically larger than the string scale. This implies that the widths of mesons should remain zero to all orders in the perturbative $\frac{1}{\sqrt{\lambda}}$ expansion.

One can also turn on a quark chemical potential $\mu < m_q^{(T)}$ in the boundary theory by setting $A_t = \mu$, where A_t is the time component of the gauge field on the D7-brane [11,12]. Since the DBI action and its higher order α' corrections contain only derivatives of A_t, the D7-embedding and the meson spectrum are not modified by turning on the constant mode of A_t. Thus, the meson widths and the net quark density remain zero to all orders in the α' expansion even at a finite chemical potential.[1]

The above conclusions can be further illuminated by simple thermodynamic reasoning. From (3), $\beta m_q^{(T)} \propto \sqrt{\lambda}$, the quark (or antiquark) density, being proportional to $e^{-\beta m_q^{(T)} \pm \beta\mu}$, is then exponentially suppressed in $\sqrt{\lambda}$ for $\mu < m_q^{(T)}$. Similarly, since the binding energy E_{BE} of a meson is $2m_q^{(T)}$, thermal effects which destabilize the mesons are also exponentially suppressed in $\sqrt{\lambda}$. Thus the meson widths and quark densities are not visible in the perturbative expansion in $1/\sqrt{\lambda}$ and can only have non-perturbative origins on the worldsheet.

There are indeed non-perturbative corrections in α' which effectively change the topology of the D7-brane embedding and generate non-vanishing meson widths and quark densities. To see this it is more convenient to analytically continue (1) to Euclidean sig-

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[1] Above $\mu > m_q^{(T)}$ a new phase, where the D7 brane falls into the horizon, is thermodynamically preferred. This phase exhibits both meson widths and finite quark density [11–14].

nature with $t \to -i\tau$. Then the r–τ plane has the topology of a disk. The angular direction τ has a period given by the inverse temperature β. The center of the disk is located at $r = r_0$. Open strings on the D7-brane are described by worldsheets with the topology of a disk whose boundary lies on the D7-brane. Denoting the worldsheet coordinates as $\rho \in [0, 1]$ and $\sigma \in [0, 2\pi)$, the worldsheets separate into different topological sectors corresponding to the winding number m of the target space disk (r, τ) wrapping around the worldsheet disk (ρ, σ). The DBI action arises from the sector of trivial winding number $m = 0$, in which (ρ, σ) is mapped to a single point on the D7-brane. In all the other (nontrivial) sectors, the string worldsheet is mapped to the region in the r–τ plane from the location of the D7-brane all the way to the horizon $r = r_0$ (see inset of Fig. 1). When analytically continued back to the Lorentzian signature, such a worldsheet describes a tunneling process in which a tiny neck is generated between the D7-brane and the black hole horizon. As a result mesons can leak through the tiny neck into the black hole and dissociate.

As an illustration, we now calculate the contributions from $m = \pm 1$ sectors to the quark density and the widths of mesons in (4). We will only be interested in the lowest order term in the α' expansion. The relevant spacetime effective action for the D7-brane can be obtained from the worldsheet path integral [15]

$$S_E[\chi^\phi] = \int\limits_{\text{disk}} DX \exp\left\{ -I[X] + \oint_{\rho=1} d\sigma\, \mu \frac{d\tau}{d\sigma} - I_B[\chi^\phi] \right\}, \qquad (6)$$

where $X = (\xi_\alpha, r, \phi)$ denotes collectively the worldsheet fields. $I[X]$ is the worldsheet action, which for our purpose can be taken to be the Nambu–Goto action for a string propagating in (1)

$$I[X] = \frac{1}{2\pi\alpha'} \int d\sigma\, d\rho\, \sqrt{\det h_{ab}} \qquad (7)$$

with $h_{ab} = G_{MN} \partial_a X^M \partial_b X^N$ the induced metric on the worldsheet and G_{MN} the Euclidean version of the metric (1). The second term in the exponential of (6) corresponds to turning on $A_\tau = -i\mu$ which gives in the Euclidean theory a nonzero (quark) chemical potential μ in the boundary YM theory. $I_B[\chi^\phi] = \oint_{\rho=1} d\sigma \frac{R^2}{2\pi\alpha'} \chi^\phi \partial_\rho \phi$ is the boundary action which couples the worldsheet to $\chi^\phi(\xi^\alpha)$. We have suppressed any dependence on spacetime and world sheet fermions. The integral (6) can be evaluated using the saddle point approximation in each topological sector [16], i.e. $S_E = S_{m=0} + S_{m=+1} + S_{m=-1} + \cdots$ with $S_{m=0} = S_{\text{DBI}}$.

For $m = \pm 1$, (7) has a classical solution given by

$$\tau = \pm \frac{\beta}{2\pi}\sigma, \qquad r = \hat{r}(\rho), \qquad \theta = 0, \qquad \phi = 0, \qquad \vec{x} = \vec{x}_0, \qquad (8)$$

where $\vec{x}_0$ is an arbitrary constant position vector and $\hat{r}(\rho)$ is chosen so that $\hat{r}(1) = r_m$ and $f\, d\tau^2 + \frac{1}{f} d\hat{r}^2 \propto d\rho^2 + \rho^2 d\sigma^2$. Eq. (8) represents the worldsheet of a string connecting the tip of the brane to the horizon with the $\pm$ sign corresponding to opposite orientations. It has a classical action $I_\pm = \beta m_q^{(T)}$, where $m_q^{(T)}$ is the effective quark mass introduced in (3). One can readily verify that (8) minimizes the action and satisfies the right boundary conditions at the D7-brane. Note that there are only three bosonic zero modes in (8),[2] since it costs energy to move away from $\theta = 0$ and the worldsheet time σ now coincides with τ. With χ^ϕ set to zero, Eq. (6) then yields

$$S_{m=\pm 1} = e^{-\beta m_q^{(T)}} e^{\pm \mu\beta} \frac{1}{g_s} D V_3, \qquad (9)$$

where the $e^{\pm \mu\beta}$ arises from the second term in the exponential of (6), V_3 is the spatial volume from integrating over the

three zero modes in (8), and the $\frac{1}{g_s}$ factor arises because we are evaluating the disk path integral. D is a real number coming from Gaussian integration around the saddle point (8) (including worldsheet fermions) whose sign we will fix from physical requirements. Identifying the Euclidean action with $\beta F(\beta, \mu)$ where $F(\beta, \mu)$ is the free energy, Eq. (9) leads to a net quark *charge* density $-\frac{2D}{g_s} e^{-\beta m_q^{(T)}} \sinh \beta\mu$, which in turn requires that D should be *negative*.[3] It is natural to interpret $S_{m=\pm 1}$ as the contributions from quarks and antiquarks separately: $S_{m=\pm 1} = -n_\pm V_3$, which from (9) leads to a quark and antiquark number density given by $n_\pm = e^{-\beta m_q^{(T)}} e^{\pm \mu\beta} \frac{1}{g_s}(-D)$. Note that $n_\pm \propto 1/g_s \propto N_c$ since quarks come in an N_c-multiplet.

In our derivation of (9) we have assumed the embedding of the D7-brane is given by that determined by the DBI action. This is justified for $\mu < m_q^{(T)}$ since the correction to the DBI action is exponentially small. When $\mu \geqslant m_q^{(T)}$, the backreactions from instantons become large and the embedding of Fig. 1 cannot be trusted.

The nonzero quark densities for nonzero $\mu < m_q^{(T)}$ have important implications for the phase structure of the theory. As discussed in [11] (see also [12]) based on the analysis of the DBI action (which corresponds to $\lambda = \infty$), at low temperature there is a phase transition in chemical potential at which the net quark charge density becomes nonzero. Our results strongly indicate this phase transition (at nonzero temperature) is smoothed to a crossover at any finite value of λ.

To find the widths of the mesons described by (4), we need to compute (6) to quadratic order in χ^ϕ. For simplicity, we will restrict to the $l = 0$ mode on the S_3. Near $\theta = 0$, the worldsheet action for ϕ is given by $\frac{R^2}{4\pi\alpha'} \int d\rho\, d\sigma\, (\partial\phi)^2$, which is free. The path integral is then straightforward to compute and yields for $S_{m=\pm 1}$

$$-\frac{R^2 n_\pm}{2\pi\alpha'\beta^2} \int d^3 x_0\, d\tau\, d\tau' \left(\chi^\phi(\tau, \vec{x}_0) \tilde{G}_D(\sigma, \sigma') \chi^\phi(\tau', \vec{x}_0) \right)_{\theta=0}, \qquad (10)$$

where $\tilde{G}_D(\sigma, \sigma') = \lim_{\rho\to 1, \rho'\to 1} \partial_\rho \partial_{\rho'} G_D(\rho, \sigma; \rho', \sigma')$ with $G_D(\rho, \sigma; \rho', \sigma')$ the Dirichlet propagator for a canonically normalized massless field on the unit disk and $\sigma = \frac{2\pi}{\beta}\tau$. Note that (10) only depends on the value of χ^ϕ at the tip of the brane and is nonlocal in the Euclidean time direction.

Treating (10) as a small perturbation to (4), one can compute the corrections to the Euclidean two-point function of the (meson) operator dual to χ^ϕ in the boundary YM theory, from which the corrections to the Lorentzian retarded function can be found by analytic continuation. One finds that the poles of the retarded function now obtain a nonzero imaginary part. Alternatively one can directly obtain the Lorentzian counterpart of (10) by analytically continuing the worldsheet disk to Lorentzian signature with $\sigma = i\eta = i2\pi t/\beta$, which gives the part of Rindler spacetime $ds^2 = d\rho^2 - \rho^2 d\eta^2$ with $\rho \leqslant 1$. The Lorentzian spacetime effective action can be obtained using the Schwinger–Keldysh contour, giving the Lorentzian equation of motion[4]

$$\partial_\alpha \left(\sqrt{-g}\, G_{\phi\phi} \partial^\alpha \chi^\phi \right) = -\frac{R^2 n_\pm \delta(\theta)}{\mu_7 \pi\alpha'\beta^2} \int dt'\, \tilde{G}_R(\eta - \eta') \chi^\phi(t'), \qquad (11)$$

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[2] In contrast, the $n = 0$ sector has eight zero modes corresponding to all directions on the D7-brane. There are also no fermionic zero modes here.

[3] There are a few other indications that D should be negative. The instanton action also induces a tadpole for the r component of the transverse fluctuations. One finds for D negative the tadpole pulls the brane toward the horizon as required by physical consistency. Also only for D negative, do the meson widths we calculated below have the correct sign.

[4] Eq. (11) can also be obtained directly from the Euclidean action (10) using the following general prescription: write down the equation of motion following from the Euclidean action; replace the Euclidean worldsheet time by the Lorentzian worldsheet time and the Euclidean propagator by the corresponding retarded propagator in the Lorentzian signature.

where $\tilde{G}_R(\eta - \eta') = \lim_{\rho \to 1, \rho' \to 1} \partial_\rho \partial_{\rho'} G_R(\rho, \eta; \rho', \eta')$ with $G_R(\rho, \eta; \rho', \eta')$ the *retarded* propagator for a massless field in the Rindler spacetime with Dirichlet boundary condition at $\rho = 1$. Fourier transforming (11) to momentum space and using $\int d\eta\, e^{i\nu\eta} \tilde{G}_R(\eta) = i\nu$, one finds that (5) is modified to

$$\hat{H}(\vec{k}, l = 0)\psi - \frac{i\omega n_\pm}{4\pi^3 \alpha' \mu_7 A} \delta(\theta)\psi(\theta) = \omega^2 \psi(\theta) \qquad (12)$$

with $A = \sqrt{-g}(-g^{tt})$. Writing the dispersion relation as $\omega = \omega_n - \frac{i}{2}\Gamma_n$ where n denotes the excitation number, and using first order perturbation theory in $n_\pm$ we find[5]

$$\Gamma_n^{(\pm 1)} = \frac{32\pi^3 \sqrt{\lambda}}{N_c m_q^2} |\psi_n(\theta = 0)|^2 n_\pm \qquad (13)$$

with $\psi_n(\theta = 0)$ eigenfunctions of (5) evaluated at the tip of the brane. Recall that $n_\pm$ are thermal densities for quarks and antiquarks and are proportional to N_c.

The ratio

$$\frac{\Gamma_n(k)}{\Gamma_n(0)} = \frac{|\psi_n(\theta = 0; \vec{k})|^2}{|\psi_n(\theta = 0; \vec{k} = 0)|^2} \qquad (14)$$

can be evaluated numerically and the results are shown in Fig. 2. For large k, the asymptotic form of the wave function, found in [8], can be used to show that the width (14) scales like k^2 for large k: $\Gamma_n(k)/\Gamma_n(0) = R_n[T/M](k/M)^2 + \mathcal{O}(k)$ for some function $R_n[T/M]$. Furthermore, for temperatures $T \ll M$ and $k \gg \frac{M^3}{T^2}$ one finds the closed form expression $\frac{\Gamma_n(k)}{\Gamma_n(0)} \approx \frac{2(4\pi)^4}{(n+2)(n+3/2)} \frac{T^4 k^2}{M^6}$. Fig. 2 has the interesting feature that the width is roughly constant for small k, but turns up quadratically around $k/M = 0.52(T_{\text{diss}}/T)^2$, which is roughly the momentum at which the meson approaches its limiting velocity $v_c(T)$. This is consistent with the conclusions based on the velocity dependence of the screened quark potential found in [4]. Note that the width as defined here is in the rest frame of the plasma, so the k^2 behavior at large k should be contrasted with the $1/k$ behavior of a relativistic decay width which comes from the usual time-dilation argument.

The plots here also share some similarities with those in [13] for momentum-dependence of meson widths obtained for $\mu > m_q^{(T)}$ with $\lambda = \infty$ where the relevant D7 brane embedding resembles a long spike reaching down to the horizon.

Our result (13) has a very simple physical interpretation as shown in the left plot of Fig. 3, in which a meson breaks apart by colliding with a quark (or antiquark) in the thermal bath. There are also processes corresponding to a meson breaking up by colliding with gluons in the thermal bath, shown in the right plot of Fig. 3. For such a process to happen the gluon should have an energy above the binding energy of the meson. The density of such gluons is thus suppressed by $e^{-2\beta m_q^{(T)}}$ and should be described by an instanton and anti-instanton so that the resulting worldsheet has trivial topology. We expect that contributions from such processes are also controlled by the value of the meson wave function at the tip of the brane, and possibly have similar growth with momentum.

Our method should be generic to any theory with a holographic dual. While the precise value of the width is clearly model-dependent, it is conceivable that the momentum dependence may apply in a wider context including QCD. In particular, our result highlights the contributions to quarkonium suppression from the
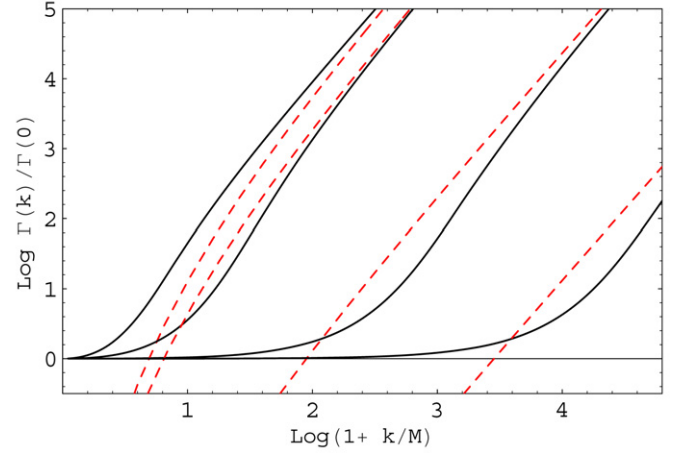

Fig. 2. The behavior of the width as a function of k for $T/T_{\text{diss}} = 0.99, 0.71, 0.3, 0.13$ from left to right. The dashed lines are analytic results for large momenta.
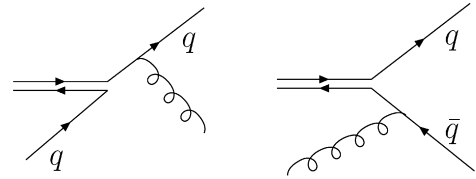

Fig. 3. *Schematic* diagrams of the relevant thermal processes contributing to the meson width. For large λ the first process is dominant, coming from the single instanton sector.

collisions with medium quarks and gluons in the large transverse momentum region in heavy ion collisions.

Consider the effect of such a momentum scaling on J/ψ with $M \approx 3$ GeV. The dissociation temperature from the gravity set-up is $T_{\text{diss}} = 0.122M \approx 370$ MeV in fairly good agreement with lattice data [17] for QCD $T_{\text{diss}} \approx 2.1T_c$ for $T_c = 170$ MeV [11]. If we take the RHIC temperature of $T = 250$ MeV (this corresponds to the second curve from the left in Fig. 2) then a moving J/Ψ's width will increase by a factor of $2(10)$ at a momentum $k = 6(13)$ GeV. When the width of a meson approaches the spacing between different meson states, the meson can no longer be considered as a well-defined quasi-particle. The momentum scaling thus implies a maximal momentum beyond which the meson no longer exists. As an illustration, suppose the width for the J/ψ in the QGP at zero momentum is about 200 MeV (which is not known) then one expects a maximal momentum around 7 GeV.

Finally, we expect the worldsheet instantons found here to have many other applications to various aspects of flavor physics in AdS/CFT.

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