Heavy quarkonium in a weakly-coupled quark-gluon plasma

Antonio Vairo

Technische Universität München

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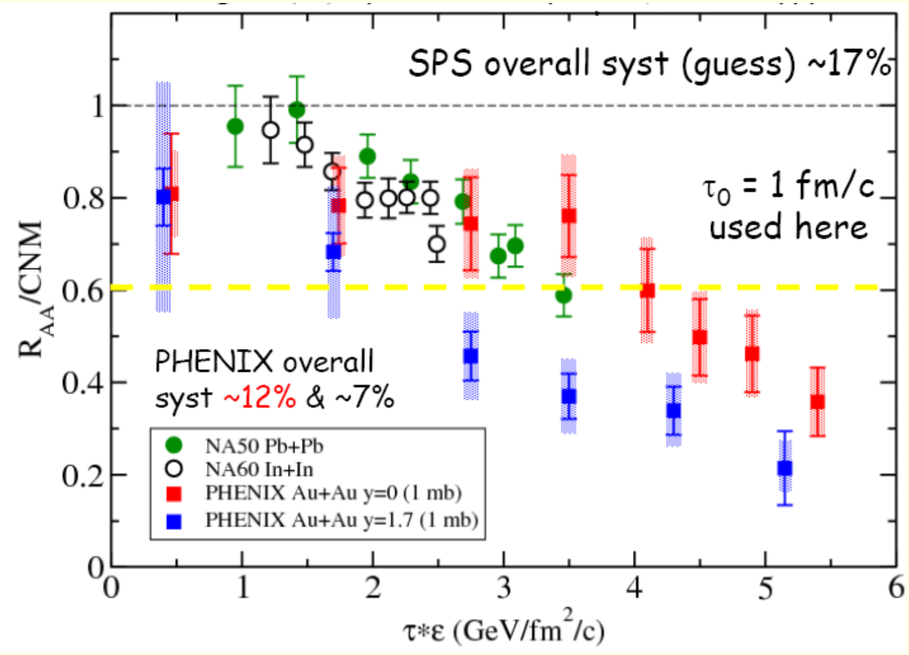
Motivation and recent literature

Original idea

- A thermal medium, like the quark-gluon plasma, induces color screening: $m_D \sim T$. If color screening induces quarkonium dissociation, then the quarkonium (dilepton rate) may provide an ideal thermometer of the quark-gluon plasma.
  - `Matsui Satz PLB 178(86)416`
- Experimental evidence of quarkonium dissociation:



  - `Leitch 08`

Recent developments

The quarkonium static potential has been calculated for $T \gg 1/r \gtrsim m_D$ by performing an analytical continuation to real time of the Euclidean Wilson loop computed in imaginary time. The calculation is done in weak-coupling resummed perturbation theory.

The imaginary part of the gluon self energy, which may be traced back to the Landau damping phenomenon, induces an imaginary part of the static potential and hence a thermal width for static quark-antiquark bound states.

- Laine Philipsen Romatschke Tassler JHEP 0703(2007)054
- Laine Philipsen Tassler JHEP 0709(2007)066
- Laine JHEP 0705(2007)028
- Burnier Laine Vepsalainen JHEP 0801(2008)043

Recent developments

- Static particles in a hot QED plasma, in real-time formalism, have been considered in the situation $1/r \sim m_D$, confirming previous results.
  ◦ `Beraudo Blaizot Ratti NPA 806(2008)312`
- A comprehensive study of non-relativistic bound states in a hot QED plasma in a non-relativistic EFT framework has been performed for a wide range of temperatures: from $T \ll m$ to $T \sim m$.

  Extrapolating to quarkonium, it was pointed out that quarkonium melts before screening sets in at a temperature

  $$\pi T_{\text{melting}} \sim m g^{4/3}$$

  ◦ `Escobedo Soto PRA 78(2008)032520, Escobedo at this workshop`

Motivation

The recent developments call for a revisitation of the physics of quarkonium states in a thermal bath in a more systematic way.

- The method:
  exploit the hierarchy of scales in the problem to construct a hierarchy of EFTs.
- The goal:
  provide a systematic description of the real-time evolution of the quarkonium state in the medium (potential, ...) and calculate its spectrum, decay width, ...

Based on

- N. Brambilla, J. Ghiglieri, P. Petreczky and A. Vairo
  Static quark-antiquark pairs at finite temperature
  Phys. Rev. D 78 (2008) 014017
- N. Brambilla, M. Escobedo, J. Ghiglieri, J. Soto and A. Vairo
  Heavy Quarkonium in a weakly-coupled quark-gluon plasma below the melting temperature
  TUM-EFT 7/10

Framework

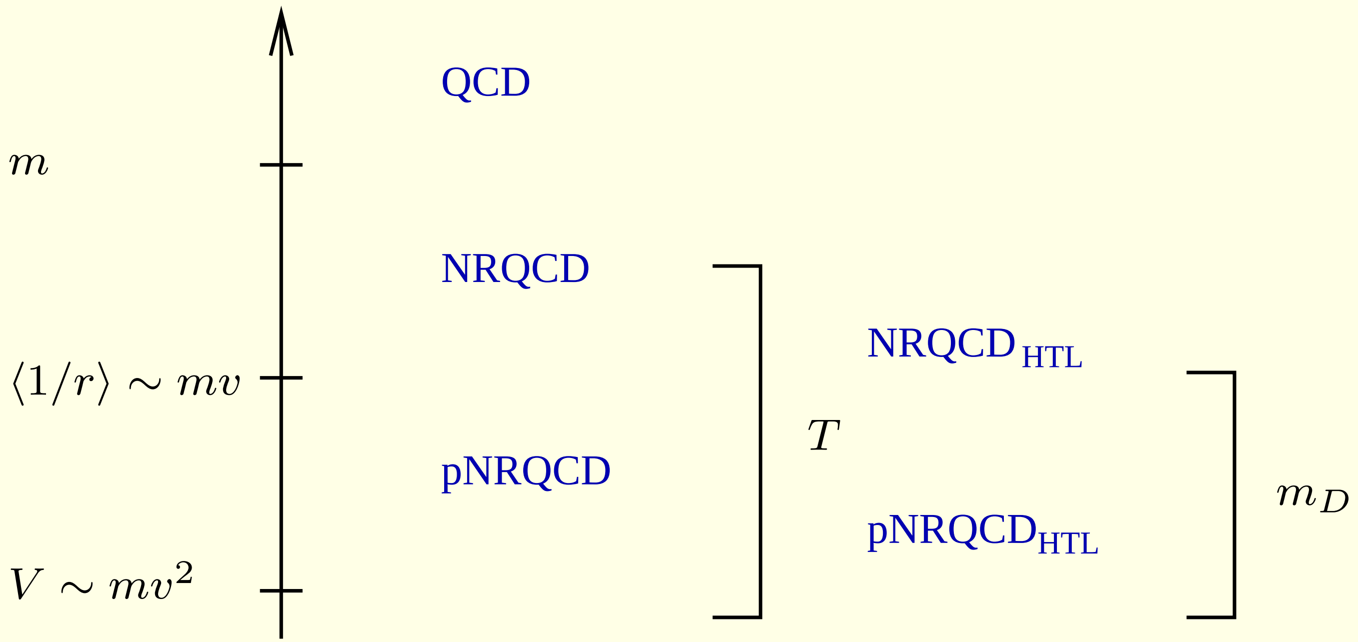
Scales

Quarkonium in a medium is characterized by different energy and momentum scales:

- the scales of the bound state (v is the relative heavy-quark velocity):
  - m (mass),
  - mv (momentum transfer, inverse distance),
  - mv^2 (kinetic energy, binding energy, potential V), ...
- the thermodynamical scales:
  - T (temperature),
  - m_D (Debye mass, i.e. screening of the chromoelectric interactions), ...

If these scales are hierarchically ordered (if the bound state is non relativistic: $v \ll 1$; in the weak coupling regime $T \gg m_D$) then we may expand physical observables in the ratio of the scales. If we separate/factorize explicitly the contributions from the different scales at the Lagrangian level this amounts to substituting QCD with a hierarchy of EFTs, which are equivalent to QCD order by order in the expansion parameters.

Effective Field Theories



∘ Brambilla Pineda Soto Vairo RMP 77(05)1423

We assume that bound states exist for

- $T \ll m$
- $\langle 1/r \rangle \sim mv \gtrsim m_D$

We neglect smaller thermodynamical scales.

Weak coupling

In the weak coupling regime:

- $v \sim \alpha_{\rm s} \ll 1$; valid for tightly bound states: $\Upsilon(1S)$, J/ψ, ...
- $T \gg gT \sim m_D$.

Effects due to the scale $\Lambda_{\rm QCD}$ will not be considered.

NRQCD

We assume $m \gg$ any other scale.

- This allows to integrate out m first and organize the EFTs as expansions in $1/m$. The first EFT is NRQCD:

$$\mathcal{L} = -\frac{1}{4} F^a_{\mu\nu} F^{a\,\mu\nu} + \sum_{i=1}^{n_f} \bar{q}_i \, i\not{D} q_i$$
$$+\psi^\dagger \left(iD_0 + \frac{\mathbf{D}^2}{2m} + \dots \right) \psi + \chi^\dagger \left(iD_0 - \frac{\mathbf{D}^2}{2m} + \dots \right) \chi + \dots$$

  ψ (χ) is the field that annihilates (creates) the (anti)fermion.

- The relevant dynamical scales of NRQCD are: $m\alpha_{\rm s}$, $m\alpha_{\rm s}^2$, ... T, m_D, ...

◦ Caswell Lepage PLB 167(1986)437
Bodwin Braaten Lepage PRD 51(1995)1125

pNRQCD

If $m\alpha_s$ is the next relevant scale, integrating it out leads to pNRQCD.

- The degrees of freedom of pNRQCD are quark-antiquark states (color singlet S, color octet O), low energy gluons and light quarks.
- The Lagrangian is organized as an expansion in $1/m$ and r:

$$
\begin{aligned}
\mathcal{L} \quad = \quad & -\frac{1}{4} F^a_{\mu\nu} F^{\mu\nu\, a} + \sum_{i=1}^{n_f} \bar{q}_i \, i \not{D} q_i \\
& + \int d^3 r \ \mathrm{Tr} \left\{ \mathrm{S}^\dagger \left(i\partial_0 - H_s \right) \mathrm{S} + \mathrm{O}^\dagger \left(iD_0 - H_o \right) \mathrm{O} \right\} \\
& \qquad + V_A \mathrm{Tr} \left\{ \mathrm{O}^\dagger \mathbf{r} \cdot g\mathbf{E} \, \mathrm{S} + \mathrm{H.c.} \right\} + \frac{V_B}{2} \mathrm{Tr} \left\{ \mathrm{O}^\dagger \mathbf{r} \cdot g\mathbf{E} \, \mathrm{O} + \mathrm{c.c.} \right\} + \cdots
\end{aligned}
$$

$$
H_s \quad = \quad \frac{\mathbf{p}^2}{m} + V_s^{(0)} + \frac{V_s^{(1)}}{m} + \dots \qquad H_o = \frac{\mathbf{p}^2}{m} + V_o^{(0)} + \frac{V_o^{(1)}}{m} + \dots
$$

○ Pineda Soto NP PC 64(1998)428

Brambilla Pineda Soto Vairo NPB 566(2000)275

pNRQCD

- At leading order in r, the singlet decouples from the octet and its EOM is:

$$(i\partial_0 - H_s)\,\mathrm{S} = 0$$

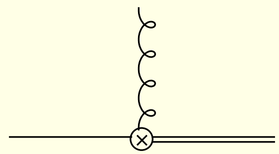
- The static potential is the Coulomb potential:

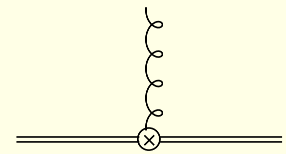
$$V_s^{(0)}(r) = -\frac{4}{3}\frac{\alpha_{\rm s}}{r}, \qquad V_o^{(0)}(r) = \frac{\alpha_{\rm s}}{6\,r}$$

- $V_A = V_B = 1 + \mathcal{O}(\alpha_{\rm s}^2)$
- Feynman rules:

 $= \theta(t)\, e^{-itH_s}$

 $= \theta(t)\, e^{-itH_o}\left(e^{-i\int dt\, A^{\rm adj}}\right)$

 $= \mathrm{O}^\dagger \mathbf{r}\cdot g\mathbf{E}\,\mathrm{S}$

 $= \mathrm{O}^\dagger \{\mathbf{r}\cdot g\mathbf{E},\, \mathrm{O}\}$

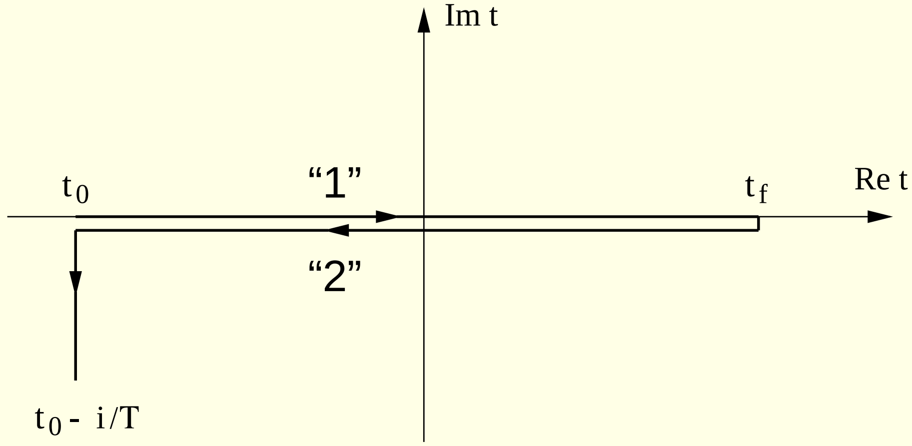
Cancellation of divergences in the spectrum I

Scale	Vacuum	Thermal
$m\alpha_{\rm s}$	$\sim m\alpha_{\rm s}^5 \dfrac{1}{\epsilon_{\rm IR}}$	—
T		
$m\alpha_{\rm s}^2$		

- Brambilla Pineda Soto Vairo PRD 60(1999)091502, PLB 470(1999)215
  Kniehl Penin NPB 563(1999)200

Real time

The contour of the partition function is modified to allow for real time:



In real time, the degrees of freedom double ("1" and "2"), however, the advantages are

- the framework becomes very close to the one for $T = 0$ EFTs;
- in the static quark sector, the second degrees of freedom, labeled "2", decouple from the physical degrees of freedom, labeled "1".

This leads to a simpler treatment with respect to alternative calculations in imaginary time formalism + analytical continuation in real time.

Quarkonium spectrum and widths below T_{melting}

Quarkonium for $m\alpha_s \gg \pi T \gg m\alpha_s^2 \gtrsim m_D$

We consider quarkonium in the following regime:

$$m\alpha_s \gg \pi T \gg m\alpha_s^2 \gtrsim m_D$$

- This implies: $mg^3 \gg T \gg mg^4$.
- πT is lower than $\pi T_{\text{melting}} \sim mg^{4/3}$: quarkonium still exists in the plasma.
- The situation may be relevant for the bottomonium $1S$ states ($\Upsilon(1S)$, η_b) at the LHC:

$$m_b\alpha_s \approx 1.5 \text{ GeV } > \pi T \approx 1 \text{ GeV } > m\alpha_s^2 \approx 0.5 \text{ GeV } \gtrsim m_D$$

- We aim at a precision in the spectrum of the order of $m\alpha_s^5$.

pNRQCD$_{\rm HTL}$

Integrating out T from pNRQCD modifies pNRQCD into pNRQCD$_{\rm HTL}$ whose

- Yang–Mills Lagrangian gets the additional hard thermal loop (HTL) part; e.g. the longitudinal gluon propagator becomes

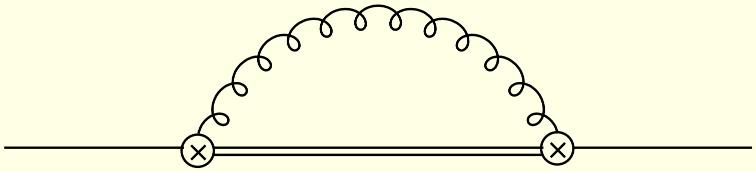
$$\frac{i}{\vec{k}^2} \to \frac{i}{k^2 + m_D^2 \left(1 - \frac{k_0}{2k} \ln \frac{k_0 + k \pm i\eta}{k_0 - k \pm i\eta}\right)}$$

  where "$+$" identifies the retarded and "$-$" the advanced propagator;

- potentials get an additional thermal correction δV to the Coulomb potential.

Integrating out T

The relevant diagram is



which reads

$$-ig^2 \, \frac{4}{3} \, \frac{r^i}{D-1} \mu^{4-D} \int \frac{d^D k}{(2\pi)^D} \, \frac{i}{E - H_o - k_0 + i\eta} \, \left[k_0^2 \, D_{ii}^{(0)}(k_0, k) + k^2 \, D_{00}^{(0)}(k_0, k) \right] r^i$$

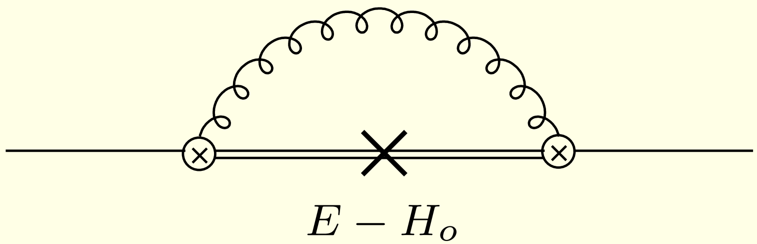
where the loop momentum region is $k_0 \sim T$ and $k \sim T$.

- Since $T \gg (E - H_o)$ we may expand

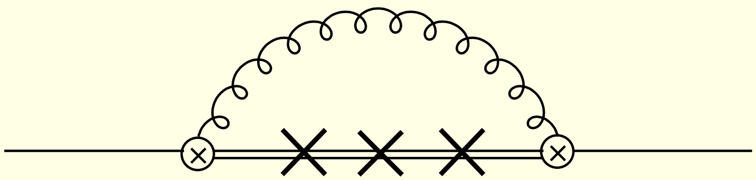
$$\frac{i}{E - H_o - k_0 + i\eta} = \frac{i}{-k_0 + i\eta} - i \frac{E - H_o}{(-k_0 + i\eta)^2} + i \frac{(E - H_o)^2}{(-k_0 + i\eta)^3} - i \frac{(E - H_o)^3}{(-k_0 + i\eta)^4} + \ldots$$

Potential: real part

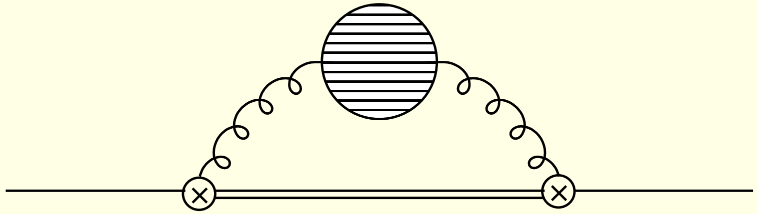
a)



b)



c)



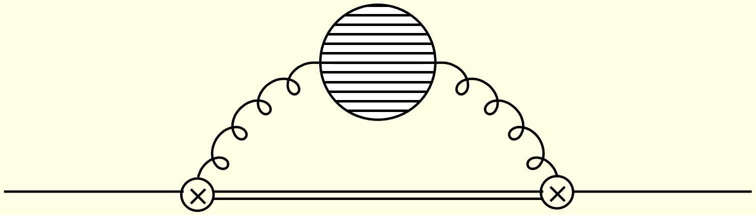
$a) \quad \sim g^2 r^2 T^3 \times E/T$

$b) \quad \sim g^2 r^2 T^3 \times (E/T)^3$

$c) \quad \sim g^2 r^2 T^3 \times (m_D/T)^2$

$$
\begin{aligned}
\mathrm{Re}\, \delta V_s(r) &= \frac{4}{9}\pi\alpha_{\mathrm{s}}^2\, r\, T^2 + \frac{8\pi}{9m}\alpha_{\mathrm{s}}\, T^2 && a) \\
&+ \frac{4\alpha_{\mathrm{s}} I_T}{3\pi}\left[-\frac{9}{8}\frac{\alpha_{\mathrm{s}}^3}{r} - \frac{17}{3}\frac{\alpha_{\mathrm{s}}^2}{mr^2} + \frac{4}{9}\frac{\pi\alpha_{\mathrm{s}}}{m^2}\delta^3(\mathbf{r}) + \frac{\alpha_{\mathrm{s}}}{m^2}\left\{\nabla_{\mathbf{r}}^2, \frac{1}{r}\right\}\right] && b) \\
&- 2\zeta(3)\frac{\alpha_{\mathrm{s}}}{\pi}\, r^2\, T\, m_D^2 + \frac{8}{3}\zeta(3)\alpha_{\mathrm{s}}^2\, r^2\, T^3 && c) \\
I_T &= \frac{2}{\epsilon} + \ln\frac{T^2}{\mu^2} - \gamma_E + \ln(4\pi) - \frac{5}{3}
\end{aligned}
$$

Potential: imaginary part



Landau-damping contribution

$$\mathrm{Im}\,\delta V_s(r) = \frac{2}{9}\alpha_s\, r^2\, T\, m_D^2 \left(-\frac{2}{\epsilon} + \gamma_E + \ln\pi - \ln\frac{T^2}{\mu^2} + \frac{2}{3} - 4\ln 2 - 2\frac{\zeta'(2)}{\zeta(2)}\right) + \frac{16\pi}{9}\ln 2\, \alpha_s^2\, r^2\, T^3$$

$$\sim g^2 r^2 T^3 \times \left(\frac{m_D}{T}\right)^2$$

- Note that $\pi T_{\mathrm{melting}} \sim m g^{4/3}$ comes from $\mathrm{Im}\,\delta V_s(r)$[Landau-damping] $\sim \alpha_s/r$.

Energy and thermal width from the scale T

$$\delta E_{n,l}^{(T)} = \frac{4\pi}{9}\,\alpha_s^2\,T^2\,\frac{a_0}{2}(3n^2 - l(l+1)) + \frac{8\pi}{9m}\alpha_s\,T^2 + \frac{E_n I_T \alpha_s^3}{3\pi}\left\{-\frac{32}{27}\frac{\delta_{l0}}{n} + \frac{200}{3}\frac{1}{n(2l+1)} - \frac{16}{3}\frac{1}{n^2} + \frac{27}{4}\right\} + \left(-2\zeta(3)\frac{\alpha_s}{\pi}\,T\,m_D^2 + \frac{8}{3}\zeta(3)\alpha_s^2\,T^3\right)\frac{a_0^2 n^2}{2}\left[5n^2 + 1 - 3l(l+1)\right]$$

$$\Gamma_{n,l}^{(T)} = \left[-\frac{2}{9}\alpha_s T m_D^2\left(-\frac{2}{\epsilon} + \gamma_E + \ln\pi - \ln\frac{T^2}{\mu^2} + \frac{2}{3} - 4\ln 2 - 2\frac{\zeta'(2)}{\zeta(2)}\right) - \frac{16\pi}{9}\ln 2\,\alpha_s^2\,T^3\right] a_0^2 n^2\left[5n^2 + 1 - 3l(l+1)\right]$$

where $E_n = -\frac{4m\alpha_s^2}{9n^2}$ and $a_0 = \frac{3}{2m\alpha_s}$

Cancellation of divergences in the spectrum II

Scale	Vacuum	Thermal
$m\alpha_{\rm s}$	$\sim m\alpha_{\rm s}^5 \dfrac{1}{\epsilon_{\rm IR}}$	
T	scaleless	$\sim -m\alpha_{\rm s}^5 \dfrac{1}{\epsilon_{\rm IR}}$
$m\alpha_{\rm s}^2$		

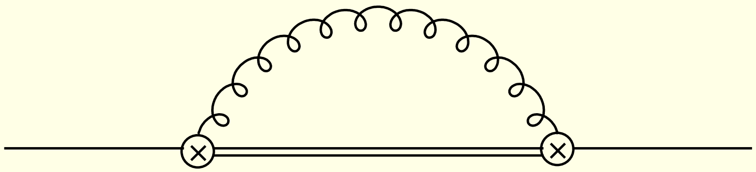
Cancellation of divergences in the spectrum II

Scale	Vacuum	Thermal
$m\alpha_s$	$\sim m\alpha_s^5 \dfrac{1}{\epsilon_{\rm IR}}$	
T	$\sim m\alpha_s^5 \left(\dfrac{1}{\epsilon_{\rm IR}} - \dfrac{1}{\epsilon_{\rm UV}}\right)$	$\sim -m\alpha_s^5 \dfrac{1}{\epsilon_{\rm IR}}$
$m\alpha_s^2$		

- Concerning the width , we anticipate that the IR divergence at the scale T will cancel against an UV divergence at the scale $m\alpha_s^2$.

Integrating out E

The relevant diagram is (gluons are HTL gluons)



where the loop momentum region is $k_0 \sim E$ and $k \sim E$.

- Since $k \sim E \ll T$ we may expand the Bose–Einstein distribution

$$n_{\rm B}(k) = \frac{T}{k} - \frac{1}{2} + \frac{k}{12\,T} + \dots\,.$$

- Since $k \sim E \gg m_D$, the HTL propagators can be expanded in $m_D^2/E^2 \ll 1$ around $D^{(0)}_{\mu\nu}$.

Energy from the scale E

$$\delta E^{E}_{n,l} = -\frac{4\pi\alpha_{\rm s}\, T m_D^2}{9}\,\frac{a_0^2 n^2}{2}\left[5n^2+1-3l(l+1)\right]$$

where $a_0 = \dfrac{3}{2m\alpha_{\rm s}}$.

- Note the complete cancellation of the vacuum contribution to the energy at the scale E (which includes the Bethe logarithm). This comes from the "$-1/2$" in the expansion of the Bose–Einstein distribution.

Cancellation of divergences in the spectrum III

Scale	Vacuum	Thermal
$m\alpha_s$	$\sim m\alpha_s^5 \frac{1}{\epsilon_{\rm IR}}$	/
T	$\sim m\alpha_s^5 \left(\frac{1}{\epsilon_{\rm IR}} - \frac{1}{\epsilon_{\rm UV}}\right)$	$\sim -m\alpha_s^5 \frac{1}{\epsilon_{\rm IR}}$
$m\alpha_s^2$	$\sim -m\alpha_s^5 \frac{1}{\epsilon_{\rm UV}}$	$\sim m\alpha_s^5 \frac{1}{\epsilon_{\rm UV}}$

Thermal width from the scale E

$$\begin{aligned}\Gamma_{n,l}^{(E)} &= 4\alpha_{\rm s}^3 T - \frac{64}{9m}\alpha_{\rm s} T E_n + \frac{32}{3}\alpha_{\rm s}^2 T \frac{1}{mn^2 a_0} \\ &+ \frac{2E_n\alpha_{\rm s}^3}{3}\left\{-\frac{32}{27}\frac{\delta_{l0}}{n} + \frac{200}{3}\frac{1}{n(2l+1)} - \frac{16}{3}\frac{1}{n^2} + \frac{27}{4}\right\} \\ &- \frac{2\alpha_{\rm s} T m_D^2}{9}\left(\frac{2}{\epsilon} + \ln\frac{E_1^2}{\mu^2} + \gamma_E - \frac{11}{3} - \ln\pi + \ln 4\right) a_0^2 n^2 [5n^2 + 1 - 3l(l+1)] \\ &+ \frac{128\alpha_{\rm s} T m_D^2}{27}\frac{\alpha_{\rm s}^2}{E_n^2} I_{n,l}\ ,\end{aligned}$$

where $E_n = -\frac{4m\alpha_{\rm s}^2}{9n^2}$ and $a_0 = \frac{3}{2m\alpha_{\rm s}}$ and $I_{1,0} = -0.49673$, $I_{2,0} = 0.64070, \ldots$

Thermal width from the scale E

$$
\begin{aligned}
\Gamma_{n,l}^{(E)} \;=\;& 4\alpha_{\rm s}^3 T - \frac{64}{9m}\alpha_{\rm s} T E_n + \frac{32}{3}\alpha_{\rm s}^2 T \frac{1}{mn^2 a_0} \\
&+ \frac{2E_n\alpha_{\rm s}^3}{3}\left\{-\frac{32}{27}\frac{\delta_{l0}}{n} + \frac{200}{3}\frac{1}{n(2l+1)} - \frac{16}{3}\frac{1}{n^2} + \frac{27}{4}\right\} \\
&- \frac{2\alpha_{\rm s} T m_D^2}{9}\left(\frac{2}{\epsilon} + \ln\frac{E_1^2}{\mu^2} + \gamma_E - \frac{11}{3} - \ln\pi + \ln 4\right) a_0^2 n^2 [5n^2 + 1 - 3l(l+1)] \\
&+ \frac{128\alpha_{\rm s} T m_D^2}{27}\frac{\alpha_{\rm s}^2}{E_n^2} I_{n,l}\ ,
\end{aligned}
$$

- The leading contribution is given by the first three terms, which are of order $\alpha_{\rm s}^3 T$. This thermal width is generated by the break up of a quark-antiquark colour-singlet state into an unbound quark-antiquark colour-octet state: a process that is kinematically allowed only in a medium. The singlet to octet break up is a different phenomenon with respect to the Landau damping. In the situation $E \gg m_D$, the first dominates over the second by a factor $(m\alpha_{\rm s}^2/m_D)^2$.

Thermal width from the scale E

$$
\begin{aligned}
\Gamma_{n,l}^{(E)} &= 4\alpha_s^3 T - \frac{64}{9m}\alpha_s T E_n + \frac{32}{3}\alpha_s^2 T \frac{1}{mn^2 a_0} \\
&+ \frac{2E_n\alpha_s^3}{3}\left\{-\frac{32}{27}\frac{\delta_{l0}}{n} + \frac{200}{3}\frac{1}{n(2l+1)} - \frac{16}{3}\frac{1}{n^2} + \frac{27}{4}\right\} \\
&- \frac{2\alpha_s T m_D^2}{9}\left(\frac{2}{\epsilon} + \ln\frac{E_1^2}{\mu^2} + \gamma_E - \frac{11}{3} - \ln\pi + \ln 4\right) a_0^2 n^2 [5n^2 + 1 - 3l(l+1)] \\
&+ \frac{128\alpha_s T m_D^2}{27}\frac{\alpha_s^2}{E_n^2} I_{n,l} \; ,
\end{aligned}
$$

- The UV divergence at the scale $m\alpha_s^2$ cancels against the IR divergence identified at the scale T.

The complete thermal contribution to the spectrum up to $\mathcal{O}(m\alpha_s^5)$

For a general number of colors N_c ($C_F = (N_c^2-1)/(2N_c)$):

$$\begin{aligned}
\delta E_{n,l}^{(\mathrm{thermal})} = {} & \frac{\pi}{9} N_c C_F\, \alpha_s^2\, T^2 \frac{a_0}{2}\left[3n^2 - l(l+1)\right] + \frac{\pi}{3} C_F^2\, \alpha_s^2\, T^2\, a_0 \\
& + \frac{E_n \alpha_s^3}{3\pi}\left[\log\left(\frac{2\pi T}{E_1}\right)^2 - 2\gamma_E\right]\left\{\frac{4C_F^3\delta_{l0}}{n} + N_c C_F^2\left[\frac{8}{n(2l+1)} - \frac{1}{n^2} - \frac{2\delta_{l0}}{n}\right]\right. \\
& \qquad \left. + \frac{2N_c^2 C_F}{n(2l+1)} + \frac{N_c^3}{4}\right\} \\
& + \frac{2E_n C_F^3 \alpha_s^3}{3\pi} L_{n,l} \\
& + \frac{a_0^2 n^2}{2}\left[5n^2 + 1 - 3l(l+1)\right]\left\{-\left[\frac{3}{2\pi}\zeta(3) + \frac{\pi}{3}\right] C_F\, \alpha_s\, T\, m_D^2\right. \\
& \qquad \left. + \frac{2}{3}\zeta(3)\, N_c C_F\, \alpha_s^2\, T^3\right\}
\end{aligned}$$

where $E_n = -\dfrac{mC_F^2\alpha_s^2}{4n^2}$, $a_0 = \dfrac{2}{mC_F\alpha_s}$ and $L_{n,l}$ is the Bethe logarithm.

The complete thermal width up to $\mathcal{O}(m\alpha_{\rm s}^5)$

For a general number of colors N_c:

$$\begin{aligned}
\Gamma_{n,l}^{(\rm thermal)} &= \frac{1}{3}N_c^2 C_F \alpha_{\rm s}^3 T + \frac{4}{3}\frac{C_F^2\alpha_{\rm s}^3 T}{n^2}(C_F + N_c) \\
&+ \frac{2E_n\alpha_{\rm s}^3}{3}\left\{\frac{4C_F^3\delta_{l0}}{n} + N_c C_F^2\left[\frac{8}{n(2l+1)} - \frac{1}{n^2} - \frac{2\delta_{l0}}{n}\right] + \frac{2N_c^2 C_F}{n(2l+1)} + \frac{N_c^3}{4}\right\} \\
&- \left[\frac{C_F}{6}\alpha_{\rm s} T m_D^2\left(\ln\frac{E_1^2}{T^2} + 2\gamma_E - 3 - \log 4 - 2\frac{\zeta'(2)}{\zeta(2)}\right) + \frac{4\pi}{9}\ln 2\; N_c C_F\,\alpha_{\rm s}^2\,T^3\right] \\
&\qquad\qquad \times\, a_0^2 n^2\left[5n^2 + 1 - 3l(l+1)\right] \\
&+ \frac{8}{3}C_F\alpha_{\rm s}\,T m_D^2\,a_0^2 n^4\,I_{n,l}
\end{aligned}$$

where $E_n = -\dfrac{mC_F^2\alpha_{\rm s}^2}{4n^2}$ and $a_0 = \dfrac{2}{mC_F\alpha_{\rm s}}$.

Conclusions I

- In a framework that makes close contact with modern effective field theories for non-relativistic bound states at zero temperature, we have studied the real-time evolution of a heavy quarkonium in a thermal bath of gluons and light quarks.
- The derived potential is neither the quark-antiquark free energy nor the internal energy. It is the real-time potential that describes the real-time evolution of a quarkonium state in a medium. It encodes all contributions coming from modes with energy and momentum larger than the binding energy.
- For $T > E$ the potential gets thermal contributions.
- Two mechanisms contribute to the thermal decay width: the imaginary part of the gluon self energy induced by the Landau damping phenomenon, and the quark-antiquark color singlet to color octet thermal break up. Parametrically, the first mechanism dominates for temperatures such that the Debye mass m_D is larger than the binding energy, while the latter dominates for temperatures such that m_D is smaller than the binding energy.
- In a medium, the quarkonium melts at a temperature $\pi T_{\rm melting} \sim m\, g^{4/3}$ before screening sets in.

Conclusions II

- We have studied in detail the situation: $m\alpha_s \gg \pi T \gg m\alpha_s^2 \gtrsim m_D$ that may be relevant for the bottomonium $1S$ states at the LHC.
- At leading order, the quarkonium masses increase quadratically with T, which implies the same functional increase in the energy of the leptons and photons produced in the electromagnetic decays.
- Electromagnetic decays occur at short distances $\sim 1/m \ll 1/T$, hence the standard NRQCD factorization formulas hold. At leading order, all the temperature dependence is encoded in the wave function at the origin. The leading temperature correction to it can be read from the potential and is of order $\sim n^4 T^2/(m^2\alpha_s)$. Hence, a quadratic dependence on the temperature should be observed in the frequency of produced leptons or photons.
- At leading order, a decay width linear with temperature is developed, which implies a tendency to decay to the continuum of colour-octet states. Hence, a consistently smaller number of vector and pseudoscalar ground states is expected to be in the sample with respect to the zero temperature case.