Quarkonium (bottomonium) in a weakly-coupled quark-gluon plasma

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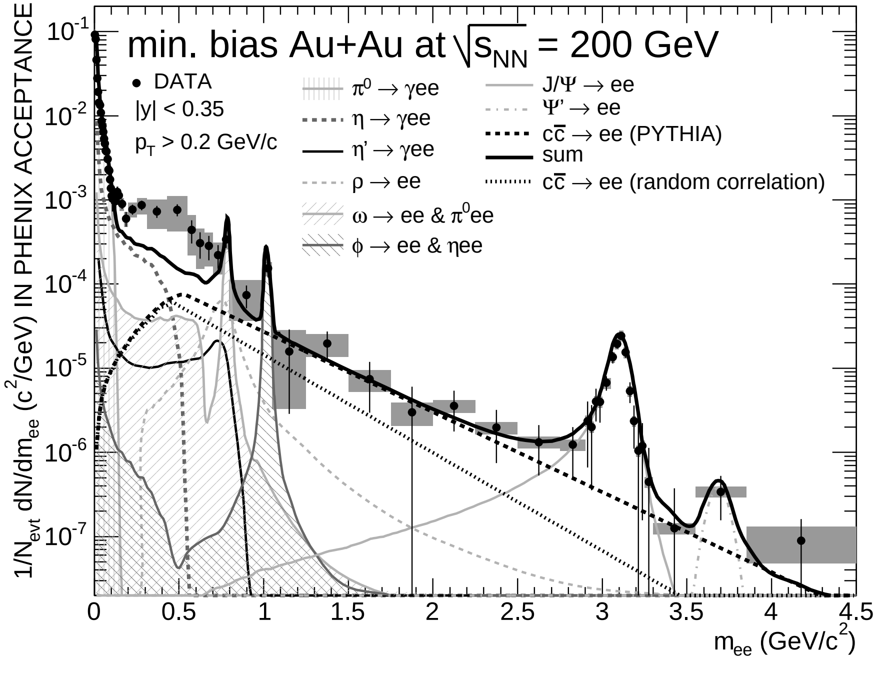
Quarkonium as a quark-gluon plasma probe

In 1986, Matsui and Satz suggest quarkonium as an ideal quark-gluon plasma probe.

- Heavy quarks are formed early in heavy-ion collisions: $1/m \sim 0.1$ fm $\ll 1$ fm.
- Heavy quarkonium formation will be sensitive to the medium.
- The dilepton signal makes the quarkonium a clean experimental probe.



Quarkonium in a thermal bath

A first step towards understanding quarkonium in a quark-gluon plasma may consist in studying how a thermal bath modifies spectrum and width of a heavy quarkonium at rest.

We will work in a well definite setting:

- weakly coupled quarkonium;
- $\pi T \gg \Lambda_{\rm QCD}$ (realized in a weakly-coupled plasma).

This situation may be realized by the bottomonium ground state at LHC.

Under these conditions, we will provide a description in terms of non-relativistic EFTs of QCD at finite temperature.

obs. What is crucial is the hierarchy of energy scales; the weakly-coupled nature allows definite perturbative calculations, but is not necessary for the EFTs set up.
Many results will hold also in a strongly coupled framework.

Scales

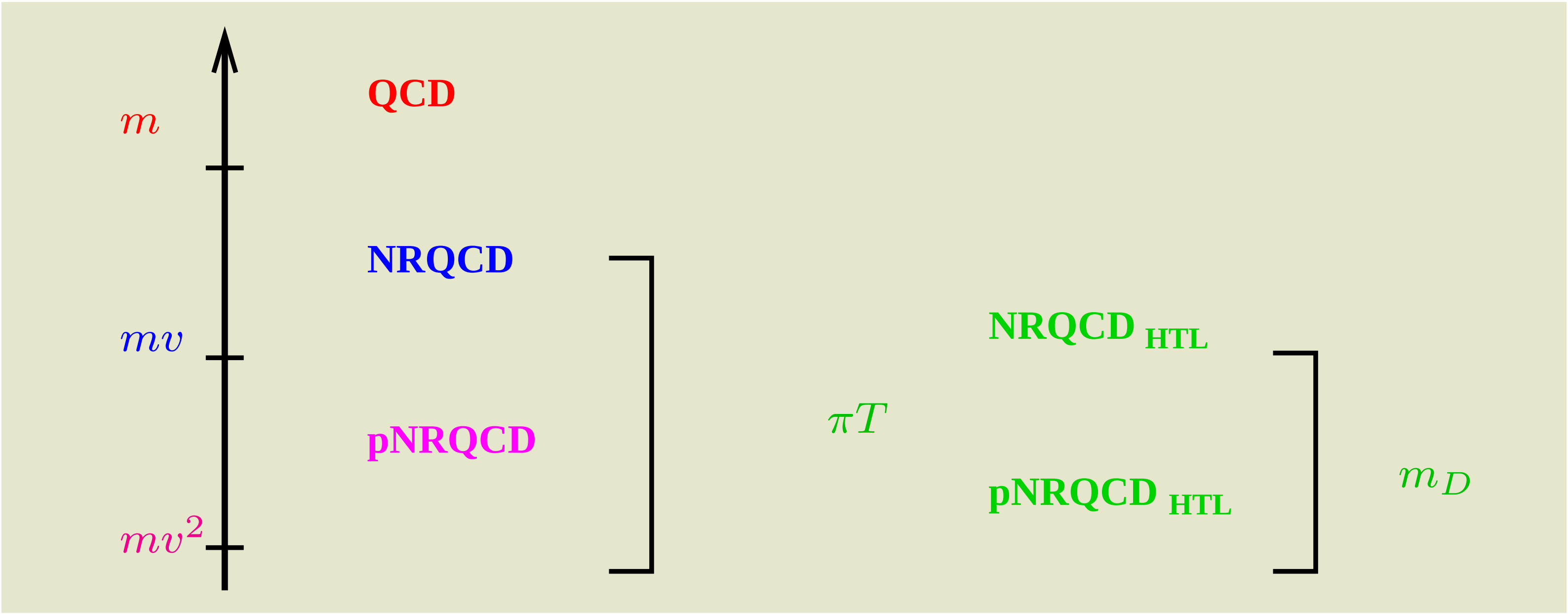
Quarkonium in a medium is characterized by different energy and momentum scales:

- the scales of a non-relativistic 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:
  - πT (temperature),
  - m_D (Debye mass, i.e. screening of the chromoelectric interactions), ...

Non-relativistic scales are hierarchically ordered: $m \gg mv \gg mv^2$,
we may assume that also the thermodynamical scales are: $\pi T \gg m_D$.

Non-relativistic Effective Field Theories at finite T



We assume that bound states exist for

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

We neglect smaller thermodynamical scales.

The bottomonium ground state at finite T

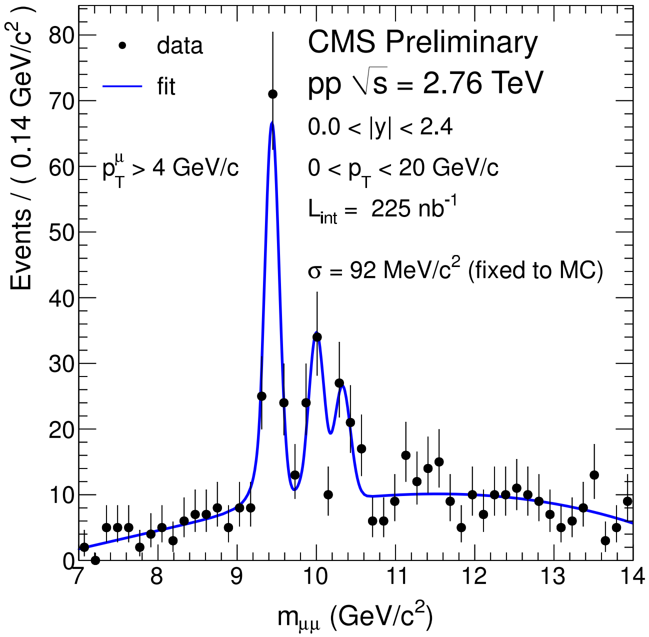
The relative size of non-relativistic and thermal scales depends on the medium and on the quarkonium state.

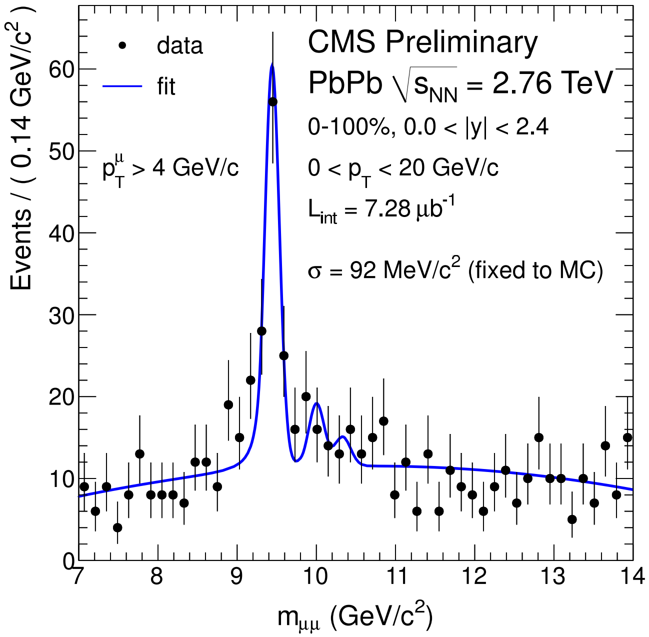
The bottomonium ground state , which is a weakly coupled non-relativistic bound state: $mv \sim m\alpha_s$, $mv^2 \sim m\alpha_s^2 \gtrsim \Lambda_{QCD}$, produced in the QCD medium of heavy-ion collisions at the LHC may possibly realize the hierarchy

$$m \approx 5 \text{ GeV} > m\alpha_s \approx 1.5 \text{ GeV} > \pi T \approx 1 \text{ GeV} > m\alpha_s^2 \approx 0.5 \text{ GeV} \gtrsim m_D, \Lambda_{QCD}$$

◦ Vairo AIP CP 1317 (2011) 241

ϒ suppression at CMS





- CMS PRL 107 (2011) 052302

NRQCD

NRQCD is obtained by integrating out modes associated with the scale m

- The Lagrangian is organized as an expansion in $1/m$:

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

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, ...

- Thermodynamical scales may be set to zero while matching.

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

pNRQCD

pNRQCD is obtained by integrating out modes associated with the scale $\frac{1}{r} \sim m\alpha_s$

- 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:

$$
\begin{aligned}
\mathcal{L}_{\rm pNRQCD} &= \int d^3r \; {\rm Tr} \left\{ {\rm S}^\dagger \left(i\partial_0 - \frac{\mathbf{p}^2}{m} + \cdots - V_s \right) {\rm S} \right. \\
&\qquad \left. + {\rm O}^\dagger \left(iD_0 - \frac{\mathbf{p}^2}{m} + \cdots - V_o \right) {\rm O} \right\} \\
&\qquad - \frac{1}{4} F^a_{\mu\nu} F^{\mu\nu\,a} + \sum_{i=1}^{n_f} \bar{q}_i \, i\not{D} q_i + \Delta\mathcal{L}
\end{aligned}
$$

$$
\Delta\mathcal{L} = \int d^3r \; V_A {\rm Tr} \left\{ {\rm O}^\dagger \mathbf{r} \cdot g\mathbf{E}\,{\rm S} + {\rm H.c.} \right\} + \frac{V_B}{2} {\rm Tr} \left\{ {\rm O}^\dagger \mathbf{r} \cdot g\mathbf{E}\,{\rm O} + {\rm c.c.} \right\} + \cdots
$$

- At leading order in r, the singlet S satisfies the Schrödinger equation.

pNRQCD

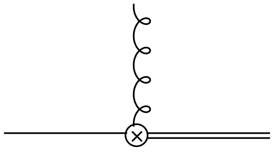
- Thermodynamical scales may be set to zero while matching.
- The static potential is the Coulomb potential:

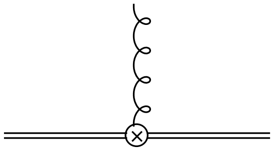
$$V_s(r) = -\frac{4}{3}\frac{\alpha_{\rm s}}{r} + \dots, \qquad\qquad V_o(r) = \frac{\alpha_{\rm s}}{6\,r} + \dots$$

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

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

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

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

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

◦ Pineda Soto NP PS 64 (1998) 428
Brambilla Pineda Soto Vairo NPB 566 (2000) 275

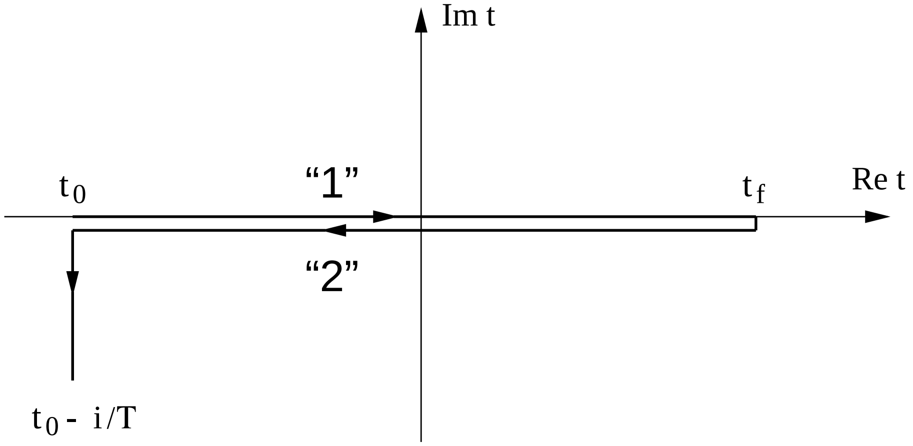
Cancellation of divergences in the spectrum I

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 heavy quark sector, the second degrees of freedom, labeled “2”, decouple from the physical degrees of freedom, labeled “1”.

Real-time gluon propagator

- Free gluon propagator in Coulomb gauge:

$$\mathbf{D}_{00}^{(0)}(k) = \mathbf{D}_{00}^{(0)}(k)_{T=0}\mathbf{1}_{2\times 2}$$

$$\mathbf{D}_{ij}^{(0)}(k) = \mathbf{D}_{ij}^{(0)}(k)_{T=0}\mathbf{1}_{2\times 2} + \left(\delta_{ij} - \frac{k^i k^j}{\mathbf{k}^2}\right) 2\pi\delta(k^2)\, n_{\mathrm{B}}(|k^0|) \begin{pmatrix} 1 & 1 \\ 1 & 1 \end{pmatrix}$$

where

$$n_{\mathrm{B}}(k^0) = \frac{1}{e^{k^0/T} - 1}$$

In Coulomb gauge, only transverse gluons carry a thermal part.

Real-time potential

- Quark-antiquark potential:

$$\mathbf{V} = \begin{pmatrix} V & 0 \\ -2i\,\mathrm{Im}\,V & -V^* \end{pmatrix}$$

- Quark-antiquark propagator:

$$\mathbf{S}^{(0)}_{\bar{Q}Q}(p) = \begin{pmatrix} \dfrac{i}{p^0 - \mathbf{p}^2/m - V + i\epsilon} & 0 \\ 2\,\mathrm{Re}\,\dfrac{i}{p^0 - \mathbf{p}^2/m - V + i\epsilon} & \dfrac{-i}{p^0 - \mathbf{p}^2/m - V^* - i\epsilon} \end{pmatrix}$$

The vanishing of the “12” component ensures that the “2” component decouples from the physical heavy quarks, i.e. the component “1”.

○ Brambilla Ghiglieri Petreczky Vairo PRD 78 (2008) 014017

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

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

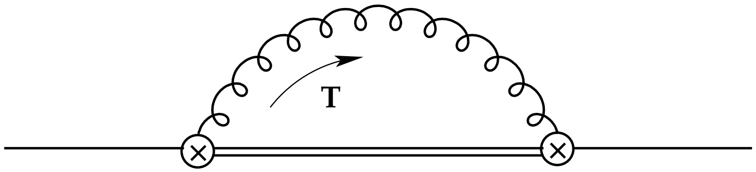
  where "$+$" identifies the retarded and "$-$" the advanced propagator;
  ○ Braaten Pisarski PRD 45 (1992) 1827

- potentials get additional thermal corrections δV.

○ Brambilla Ghiglieri Petreczky Vairo PRD 78 (2008) 014017
Escobedo Soto PRA 78 (2008) 032520
Vairo PoS CONFINEMENT8 (2008) 002

Integrating out T

The relevant diagram is



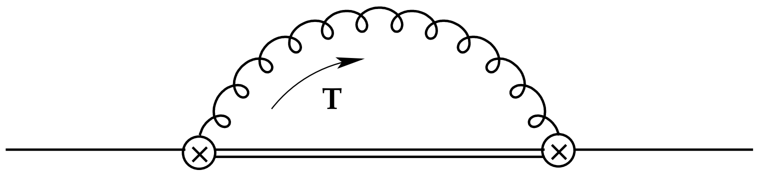
and radiative corrections. The loop momentum region is $k_0 \sim T$ and $k \sim T$.

- Since $T \gg (E - \mathbf{p}^2/m - V_o)$ we may expand the octet propagator

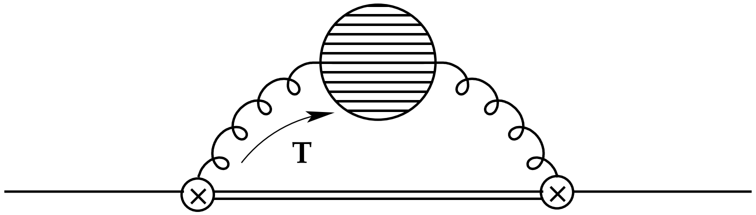
$$\frac{i}{E - \mathbf{p}^2/m - V_o - k_0 + i\eta} = \frac{i}{-k_0 + i\eta} - i\frac{E - \mathbf{p}^2/m - V_o}{(-k_0 + i\eta)^2} + \dots$$

Integrating out T: real potential

a)



b)

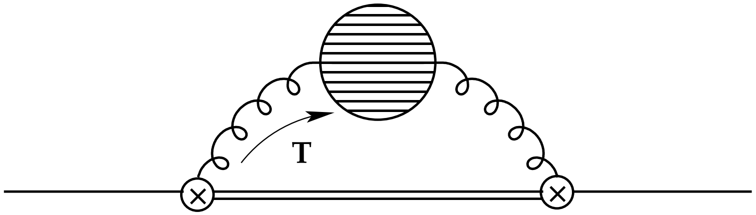


$\sim g^2 r^2 T^3 \times E/T$

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

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

Integrating out T: imaginary potential



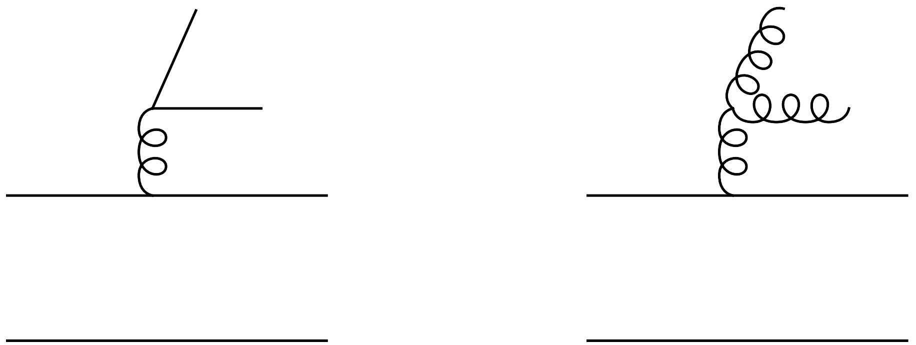
Landau-damping contribution

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

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

Landau damping

The Landau damping phenomenon originates from the scattering of the quarkonium with hard space-like particles in the medium.



○ `Laine Philipsen Romatschke Tassler JHEP 0703 (2007) 054`

- When $\mathrm{Im}\, V_s(r)|_{\mathrm{Landau-damping}} \sim \mathrm{Re}\, V_s(r) \sim \alpha_{\mathrm{s}}/r$, the quarkonium dissociates:

$$\pi T_{\mathrm{dissociation}} \sim mg^{4/3}$$

- When $\langle 1/r \rangle \sim m_D$, the interaction is screened; note that

$$\pi T_{\mathrm{screening}} \sim mg \gg \pi T_{\mathrm{dissociation}}$$

$\Upsilon(1S)$ dissociation temperature

The $\Upsilon(1S)$ dissociation temperature:

m_c (MeV)	$T_{\text{dissociation}}$ (MeV)
∞	480
5000	480
2500	460
1200	440
0	420

A temperature πT about 1 GeV is below the dissociation temperature.

- Escobedo Soto PRA 82 (2010) 042506

Energy and thermal width from the scale T

$$\delta E_{1S}^{(T)} = \frac{2\pi}{3}\,\alpha_s^2\,T^2 a_0 + \frac{8\pi}{9m}\alpha_s\,T^2 + \frac{7225}{324}\,\frac{E_1 I_T \alpha_s^3}{\pi} - \zeta(3)\alpha_s\,T\left(6\frac{m_D^2}{\pi} - 8\alpha_s\,T^2\right)a_0^2$$

$$\Gamma_{1S}^{(T)} = \left[-\frac{4}{3}\alpha_s T m_D^2\left(-\frac{2}{\epsilon} + \gamma_E + \ln\pi - \ln\frac{T^2}{\mu^2} + \frac{2}{3} - 4\ln 2 - 2\frac{\zeta'(2)}{\zeta(2)}\right) - \frac{32\pi}{3}\ln 2\,\alpha_s^2\,T^3\right]a_0^2$$

where $E_1 = -\frac{4m\alpha_s^2}{9}$ and $a_0 = \frac{3}{2m\alpha_s}$

- Brambilla Escobedo Ghiglieri Soto Vairo JHEP 1009 (2010) 038

Cancellation of divergences in the spectrum II

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

Cancellation of divergences in the spectrum II

Scale	Vacuum	Thermal
$m\alpha_{\mathrm{s}}$	$\sim m\alpha_{\mathrm{s}}^5 \frac{1}{\epsilon_{\mathrm{IR}}}$	
T	$\sim m\alpha_{\mathrm{s}}^5 \left(\frac{1}{\epsilon_{\mathrm{IR}}} - \frac{1}{\epsilon_{\mathrm{UV}}}\right)$	$\sim -m\alpha_{\mathrm{s}}^5 \frac{1}{\epsilon_{\mathrm{IR}}}$
$m\alpha_{\mathrm{s}}^2$		

- Concerning the width , the IR divergence at the scale T will cancel against an UV divergence at the scale $m\alpha_{\mathrm{s}}^2$.

Quasi-free dissociation

For general T, the thermal width due to Landau damping reads

$$\Gamma_{1S} = \sum_p \int_{q_{\rm min}} \frac{d^3q}{(2\pi)^3}\, f_p(|\mathbf{q}|)(1 \pm f_p(|\mathbf{q}|))\, \sigma_{1S}(|\mathbf{q}|)\,,$$

where the sum runs over the different incoming light partons and $f_g = n_{\rm B}$ or $f_q = n_{\rm F}$.

σ_{1S} is known as the quarkonium quasi-free dissociation cross section.
The thermal NR EFTs provide analytic expressions of σ_{1S} for different temperatures.

○ Brambilla Escobedo Ghiglieri Vairo TUM-EFT 27/11

Quasi-free dissociation: previous literature

In the previous literature, it was assumed

$$\Gamma_{1S} = \sum_p \int_{q_{\min}} \frac{d^3q}{(2\pi)^3} \, f_p(|\mathbf{q}|) \, \sigma_{\rm HQ}(|\mathbf{q}|) \, ,$$

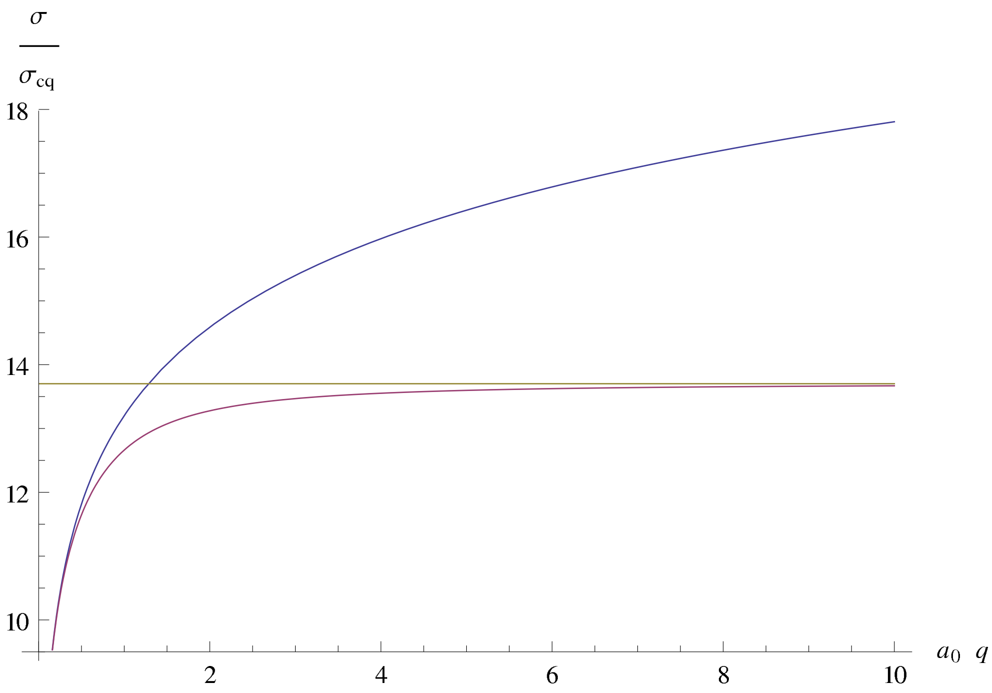
with $\sigma_{\rm HQ} = 2\sigma_c$, where σ_c is the cross section for the process $pc \to pc$ at $T = 0$.

∘ `Grandchamp Rapp, PLB 523 (2001) 60, ...`

The EFT analysis proves this assumption to be incorrect, because

- the dependence on the thermal distributions of the incoming and outgoing partons is different;
- σ_{1S} cannot be identified with $\sigma_{\rm HQ}$, moreover it is temperature dependent.

Quasi-free dissociation: light-quark contribution



$m_D a_0 = 0.001$

blue line: $mv \gg T \gg m_D \gg E$
(dipole approximation)

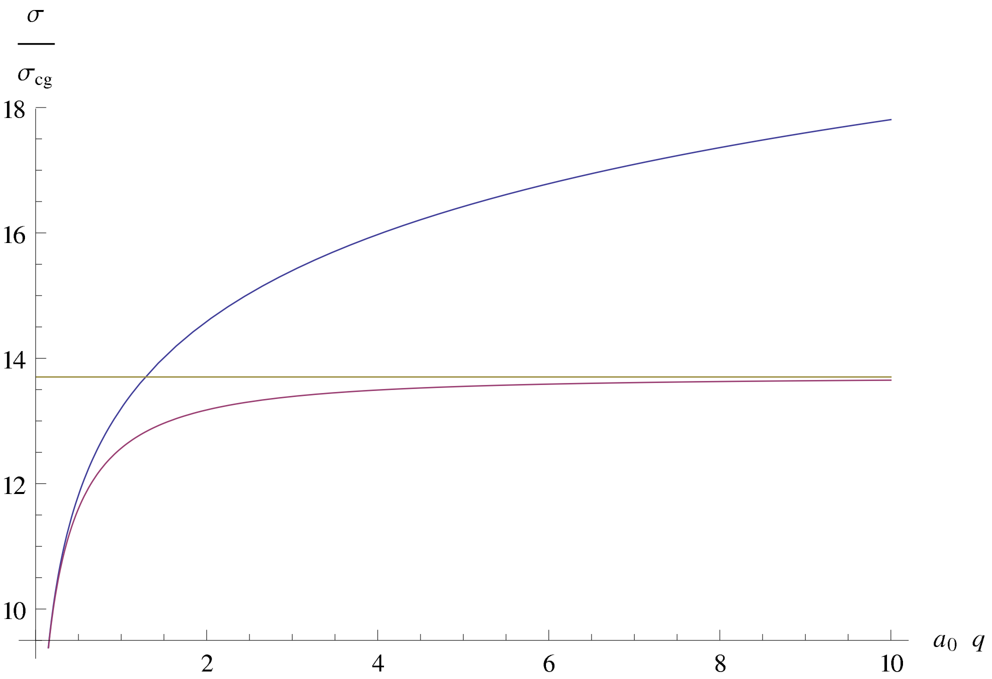
pink line: $T \sim mv \gg m_D$

yellow line: $T \gg mv \sim m_D$

$\sigma_{cq} \equiv 8\pi C_F n_f \, \alpha_s^2 \, a_0^2$

- Brambilla Escobedo Ghiglieri Vairo TUM-EFT 27/11

Quasi-free dissociation: gluon contribution



$m_D a_0 = 0.001$

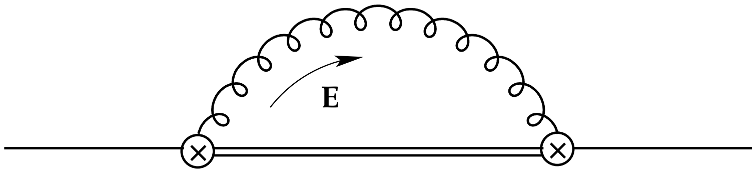
blue line: $mv \gg T \gg m_D \gg E$
(dipole approximation)
pink line: $T \sim mv \gg m_D$
yellow line: $T \gg mv \sim m_D$

$\sigma_{cg} \equiv 8\pi C_F N_c \, \alpha_s^2 \, a_0^2$

∘ Brambilla Escobedo Ghiglieri Vairo TUM-EFT 27/11

Integrating out E

The relevant diagram is



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

- Gluons are HTL gluons.
- 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 $m_D^2/E^2 \ll 1$.

Momentum regions

In the loop with transverse gluons, this type of integral appears

$$\int \frac{d^{D-1}k}{(2\pi)^{D-1}} \int_0^\infty \frac{dk_0}{2\pi} \frac{1}{k_0^2 - k^2 - m_D^2 + i\eta} \left(\frac{1}{E - H_o - k_0 + i\eta} + \frac{1}{E - H_o + k_0 + i\eta} \right)$$

which exhibits two momentum regions

- off-shell region: $k_0 - k \sim E$, $\quad k_0 \sim E$, $\quad k \sim E$;
- collinear region: $k_0 - k \sim m_D^2/E$, $\quad k_0 \sim E$, $\quad k \sim E$.

In our energy scale hierarchy, the collinear scale is $mg^4 \gg m_D^2/E \gg mg^6$, i.e. it is smaller than m_D by a factor of $m_D/E \ll 1$ and still larger than the non-perturbative magnetic mass, which is of order g^2T, by a factor $T/E \gg 1$.

Integrating out E: energy

$$\delta E_{1S}^{E} = -\frac{4\pi\alpha_{\rm s}\,T m_D^2}{3} a_0^2$$

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.

○ `Brambilla Escobedo Ghiglieri Soto Vairo JHEP 1009 (2010) 038`

Cancellation of divergences in the spectrum III

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

Integrating out E: thermal width

$$
\begin{aligned}
\Gamma_{1S}^{(E)} \quad = \quad & 4\alpha_s^3 T - \frac{64}{9m}\alpha_s T E_1 + \frac{32}{3}\alpha_s^2 T \frac{1}{m a_0} + \frac{7225}{162} E_1 \alpha_s^3 \\
& - \frac{4\alpha_s T m_D^2}{3}\left(\frac{2}{\epsilon} + \ln\frac{E_1^2}{\mu^2} + \gamma_E - \frac{11}{3} - \ln\pi + \ln 4\right) a_0^2 \\
& + \frac{128\alpha_s T m_D^2}{81}\frac{\alpha_s^2}{E_1^2} I_{1,0}
\end{aligned}
$$

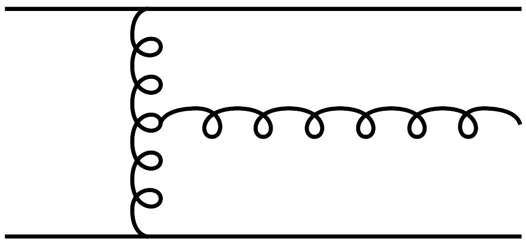
where $E_1 = -\dfrac{4m\alpha_s^2}{9}$ and $a_0 = \dfrac{3}{2m\alpha_s}$ and $I_{1,0} = -0.49673$ (similar to the Bethe log).

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

◦ Brambilla Escobedo Ghiglieri Soto Vairo JHEP 1009 (2010) 038

Singlet to octet break up

The thermal width at the scale E, which is of order $\alpha_s^3 T$, is generated by the break up of a quark-antiquark color-singlet state into an unbound quark-antiquark color-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, the relative size of which is $(E/m_D)^2$. In the situation $m\alpha_s^2 \gg m_D$, the first dominates over the second by a factor $(m\alpha_s^2/m_D)^2$.

○ Brambilla Ghiglieri Petreczky Vairo PRD 78 (2008) 014017

Gludissociation

For general T, the thermal width due to $\mathrm{S} \to \mathrm{O} + g$ break up in a medium reads

$$\Gamma_{1S} = \int_{|\mathbf{q}|\geq|E_{1S}|} \frac{d^3q}{(2\pi)^3}\, n_{\mathrm{B}}(|\mathbf{q}|)\, \sigma_{1S}(|\mathbf{q}|) \xrightarrow[T\gg E]{} \text{previous slide}$$

with

$$\sigma_{1S}(|\mathbf{q}|) = \frac{\alpha_{\mathrm{s}} C_F}{3} 2^{10}\pi^2\rho(\rho+2)^2 \frac{E_1^4}{m|\mathbf{q}|^5}\left(t(|\mathbf{q}|)^2+\rho^2\right) \frac{\exp\left(\frac{4\rho}{t(|\mathbf{q}|)}\arctan\left(t(|\mathbf{q}|)\right)\right)}{e^{\frac{2\pi\rho}{t(|\mathbf{q}|)}}-1}$$

where $\rho \equiv 1/(N_c^2-1)$ and $t(|\mathbf{q}|) \equiv \sqrt{|\mathbf{q}|/|E_1|-1}$.

σ_{1S}, which is the cross section of the process $\mathrm{S} \to \mathrm{O} + g$ in the vacuum, is known as the quarkonium gludissociation cross section.

- Brambilla Escobedo Ghiglieri Vairo JHEP 1112 (2011) 116
  Brezinski Wolschin PLB 707 (2012) 534

Gluodissociation: the Bhanot–Peskin approximation

In the large N_c limit:

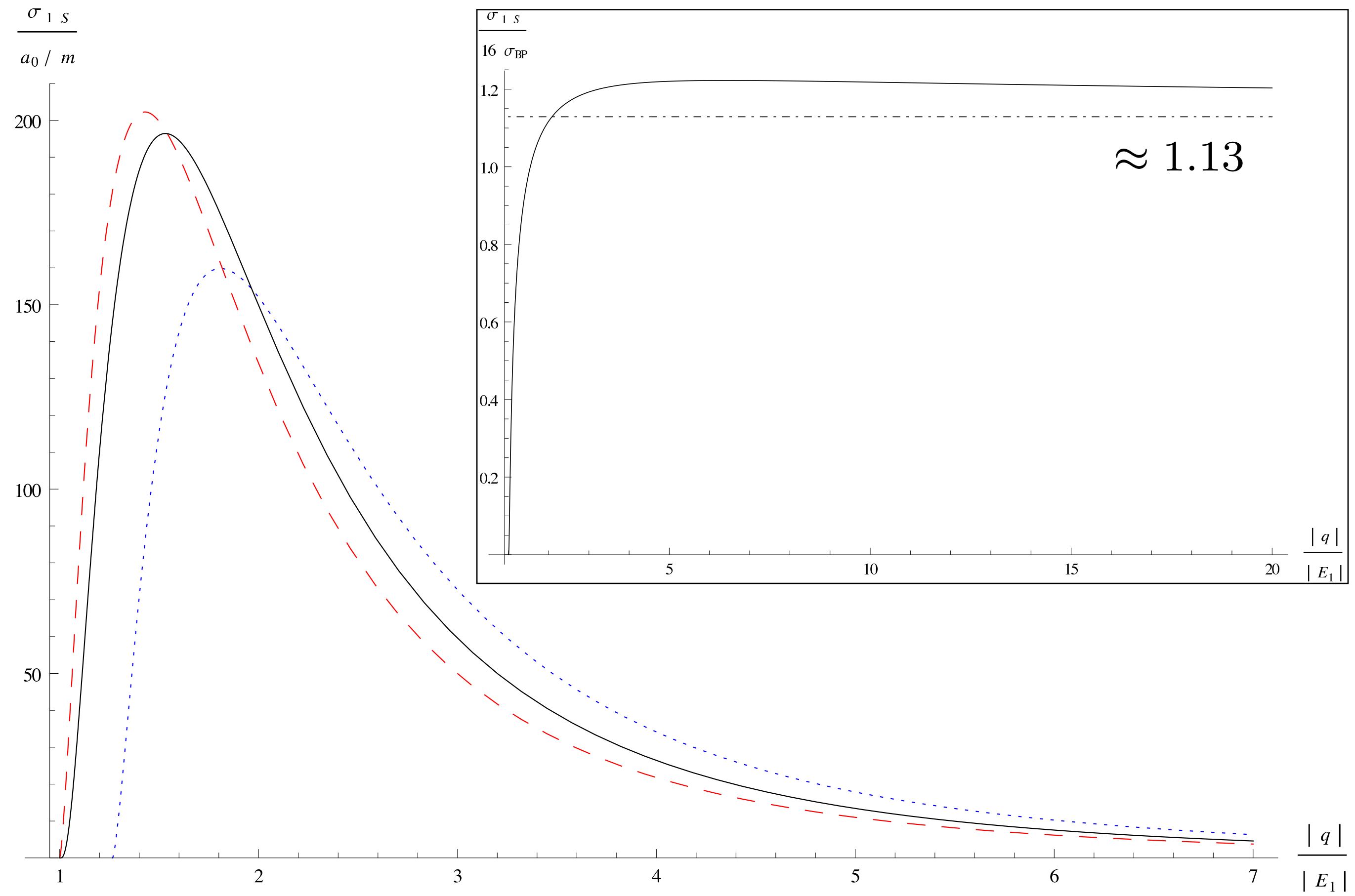
$$\sigma_{1S}(|\mathbf{q}|) \xrightarrow[N_c\to\infty]{} 16\,\frac{2^9\pi\alpha_{\rm s}}{9}\,\frac{|E_1|^{5/2}}{m}\,\frac{(|\mathbf{q}|+E_1)^{3/2}}{|\mathbf{q}|^5} = 16\,\sigma_{1S,\rm BP}(|\mathbf{q}|)$$

$$\Gamma_{1S} \xrightarrow[N_c\to\infty]{} \int_{|\mathbf{q}|\geq|E_1|}\frac{d^3q}{(2\pi)^3}\,n_{\rm B}(|\mathbf{q}|)\,16\,\sigma_{1S,\rm BP}(|\mathbf{q}|) = \Gamma_{1S,\rm BP}$$

- Bhanot Peskin NPB 156 (1979) 391

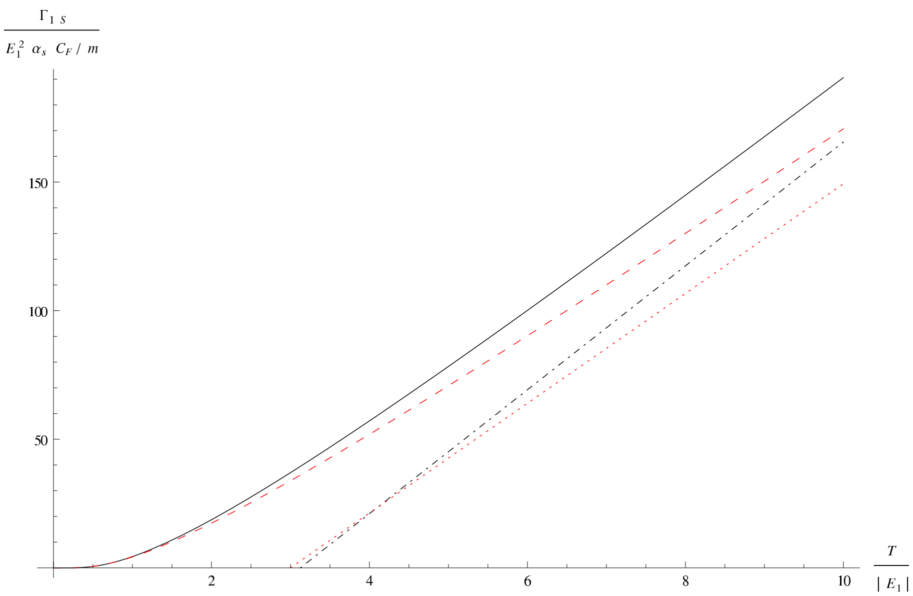
The Bhanot–Peskin (BP) approximation corresponds to neglecting final state interactions, i.e. the rescattering of a $Q\bar{Q}$ pair in a color octet configuration.

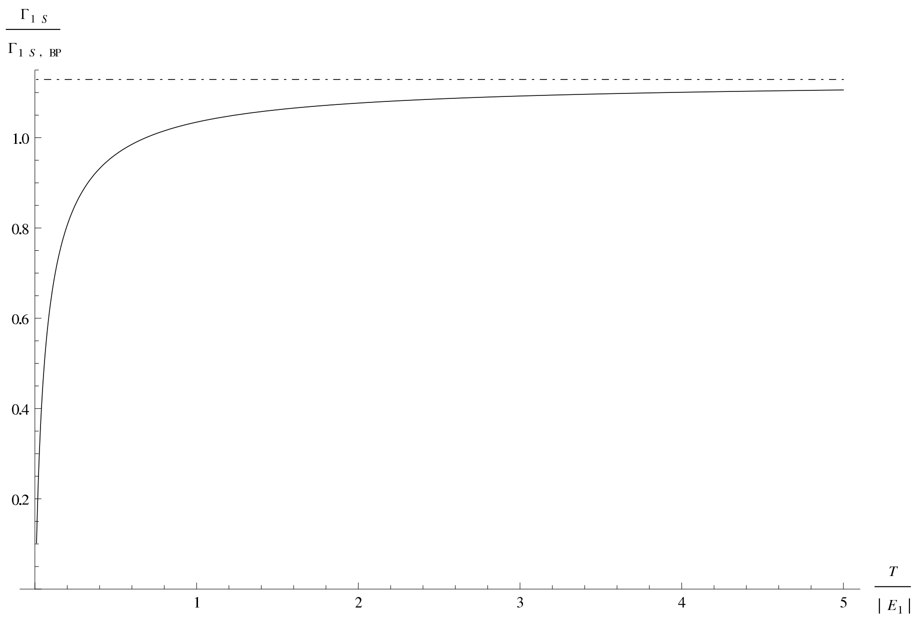
Gluodissociation: full cross section vs BP cross section



- Brambilla Escobedo Ghiglieri Vairo JHEP 1112 (2011) 116

Gluodissociation: full width vs BP width



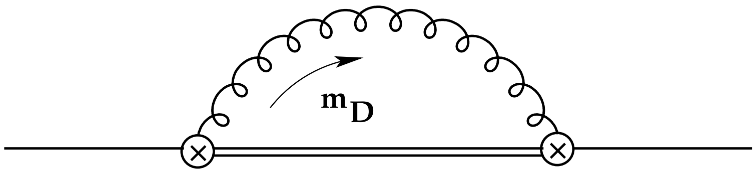


Lines on the left correspond to the $T \gg |E_1|$ analytic results of the previous slides.

- Brambilla Escobedo Ghiglieri Vairo JHEP 1112 (2011) 116

Integrating out m_D

The relevant diagram is (gluons are HTL gluons)



where the loop momentum region is $k_0 \sim m_D$ and $k \sim m_D$. This contribution is negligible with respect to the other terms calculated.

The complete mass and width up to $\mathcal{O}(m\alpha_s^5)$

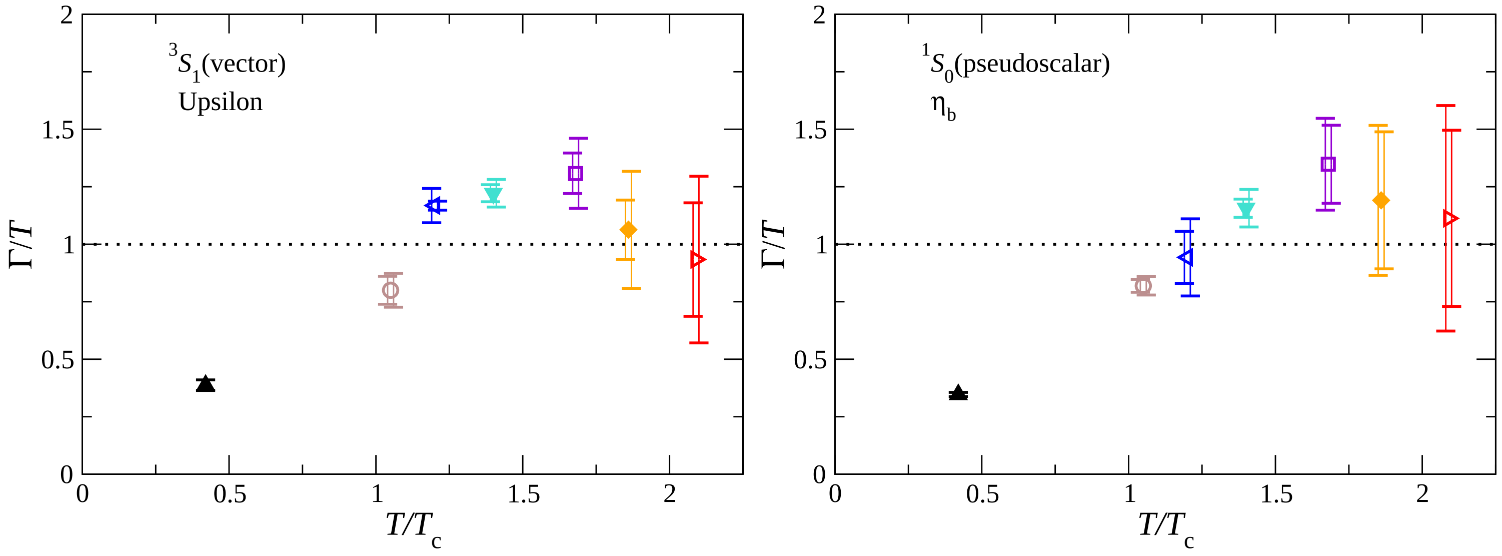
$$\delta E_{1S}^{(\mathrm{thermal})} = \frac{34\pi}{27}\alpha_s^2\, T^2 a_0 + \frac{7225}{324}\,\frac{E_1\alpha_s^3}{\pi}\left[\ln\left(\frac{2\pi T}{E_1}\right)^2 - 2\gamma_E\right] + \frac{128 E_1\alpha_s^3}{81\pi} L_{1,0} - 3a_0^2\left\{\left[\frac{6}{\pi}\zeta(3) + \frac{4\pi}{3}\right]\alpha_s\, T\, m_D^2 - \frac{8}{3}\zeta(3)\,\alpha_s^2\, T^3\right\}$$

$$\Gamma_{1S}^{(\mathrm{thermal})} = \frac{1156}{81}\alpha_s^3 T + \frac{7225}{162} E_1\alpha_s^3 + \frac{32}{9}\alpha_s\, T m_D^2\, a_0^2\, I_{1,0} - \left[\frac{4}{3}\alpha_s T m_D^2\left(\ln\frac{E_1^2}{T^2} + 2\gamma_E - 3 - \ln 4 - 2\frac{\zeta'(2)}{\zeta(2)}\right) + \frac{32\pi}{3}\ln 2\,\alpha_s^2\, T^3\right] a_0^2$$

where $E_1 = -\dfrac{4m\alpha_s^2}{9}$, $a_0 = \dfrac{3}{2m\alpha_s}$ and $L_{1,0}$ (similar $I_{1,0}$) is the Bethe logarithm.

- Brambilla Escobedo Ghiglieri Soto Vairo JHEP 1009 (2010) 038

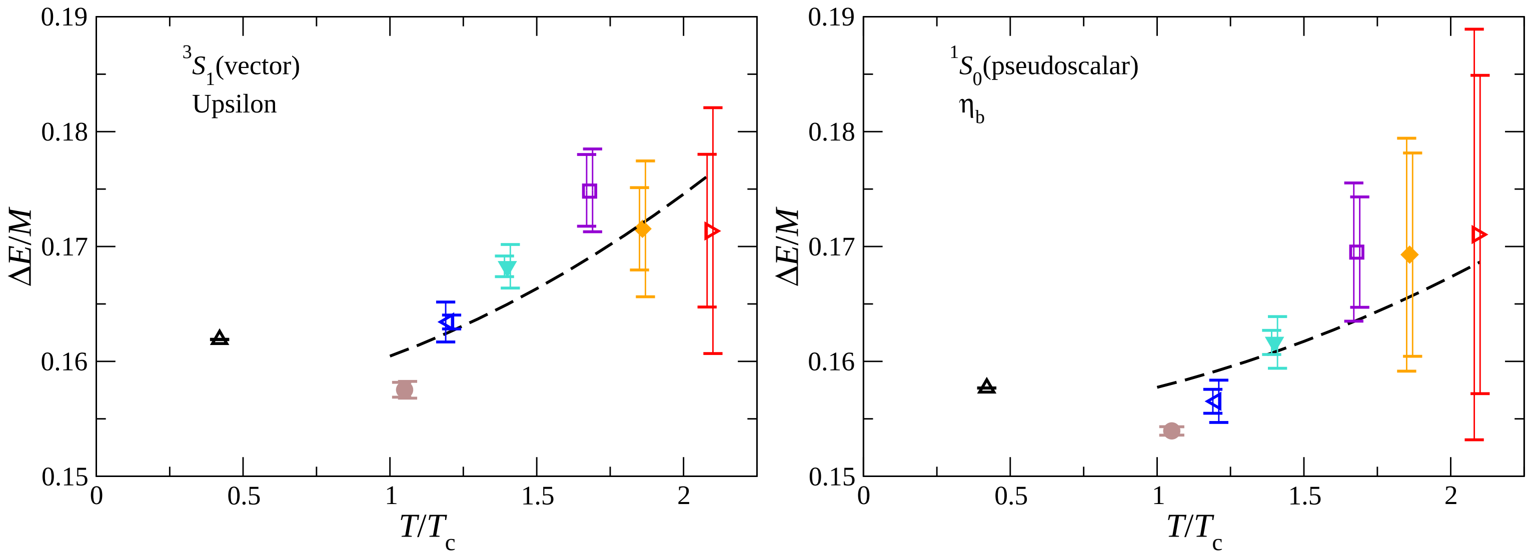
Lattice width



Consistent with $$\Gamma_{1S}^{(\text{thermal})} = \frac{1156}{81}\alpha_s^3 T \qquad \Rightarrow \alpha_s \approx 0.4.$$

- Aarts Allton Kim Lombardo Oktay Ryan Sinclair Skullerud JHEP 1111 (2011) 103

Lattice energy



Consistent with $\delta E_{1S}^{(\text{thermal})} = \frac{17\pi}{9}\alpha_s \frac{T^2}{m}$ using $\alpha_s = 0.4$ and $m = 5$ GeV.

- Aarts Allton Kim Lombardo Oktay Ryan Sinclair Skullerud
  JHEP 1111 (2011) 103

Conclusions

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 (specifically the $\Upsilon(1S)$) in a thermal bath of gluons and light quarks.

- For $T < E$ the potential coincides with the $T = 0$ one.

- 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 (aka quasi-free dissociation), and the quark-antiquark color singlet to color octet thermal break up (aka gluodissociation).

- Quarkonium dissociates at a temperature $\pi T_{\mathrm{dissociation}}$, which may be lower than the screening one.