Quarkonium suppression

an EFT approach to open quantum systems

Antonio Vairo

Technische Universität München

Based on

(1) N. Brambilla, M.A. Escobedo, J. Soto and A. Vairo
Heavy quarkonium suppression in a fireball
Phys. Rev. **D97** (2018) 074009 `arXiv:1711.04515`

(2) N. Brambilla, M.A. Escobedo, J. Soto and A. Vairo
Quarkonium suppression in heavy-ion collisions: an open quantum system approach
Phys. Rev. **D96** (2017) 034021 `arXiv:1612.07248`

Open Quantum Systems

Density matrix

An arbitrary statistical ensemble of quantum states can be represented by a density matrix ρ, which is

- Hermitian: $\rho^\dagger = \rho$;
- positive: $\langle\psi|\rho|\psi\rangle \geq 0$ for all nonzero states $|\psi\rangle$;
- and can be normalized to have unit trace: $\mathrm{Tr}\{\rho\} = 1$.

The time evolution of the density matrix is described by the von Neumann equation:

$$i\frac{d\rho}{dt} = [H, \rho]$$

which follows from the Schrödinger equation for $|\psi\rangle$. The evolution equation

- is linear in ρ;
- preserves the trace of ρ;
- is Markovian.

Open quantum system

In quantum information theory, one separates the full system into a subsystem of interest and its environment. A density matrix ρ for the subsystem can be obtained from the density matrix ρ_{full} for the full system by the partial trace over the environment states:

$$\rho = \mathrm{Tr}_{\mathrm{environment}} \{\rho_{\mathrm{full}}\}$$

In general the evolution of ρ is non-Markovian.

The evolution is Markovian if the time during which the subsystem is observed is much larger than the time scale for correlations between the subsystem and the environment. We must also restrict to the low-frequency behavior of the subsystem, which can be accomplished by smoothing out over times larger than the correlation time scale.

Lindblad equation

The density matrix ρ for the subsystem necessarily satisfies the three basic properties: it is Hermitian, positive, and it can be normalized.

If further the time evolution is linear in ρ, preserves the trace of ρ, is Markovian and the linear operator that determines the time evolution of ρ is completely positive
$\Rightarrow$ then this require the time evolution equation to have the Lindblad form

$$\frac{d\rho}{dt} = -i[H,\rho] + \sum_i (C_i \rho C_i^\dagger - \frac{1}{2}\{C_i^\dagger C_i, \rho\})$$

where H is a Hermitian operator and the C_n's are an additional set of operators called collapse operators.

- Lindblad CMP 48 (1976) 119
  Gorini Kossakowski Sudarshan JMP 17 (1976) 821

Numerical solutions: QuTiP

There exist numerical toolboxes for open quantum systems. An example is QuTiP.

QuTiP

Quantum Toolbox in Python

- Johansson, Nation, Nori, CPC 183 (2012) 1760, 184 (2013) 1234

Quarkonium in medium
as an open quantum system

Quarkonium as a multiscale system

Quarkonium being a composite system it is characterized by several energy scales:

- the scales of a non-relativistic bound state
  (v is the relative heavy-quark velocity; $v \sim \alpha_{\rm s}$ for a Coulombic bound state):
  M (mass),
  Mv (momentum transfer, inverse distance),
  Mv^2 (kinetic energy, binding energy, potential V), ...

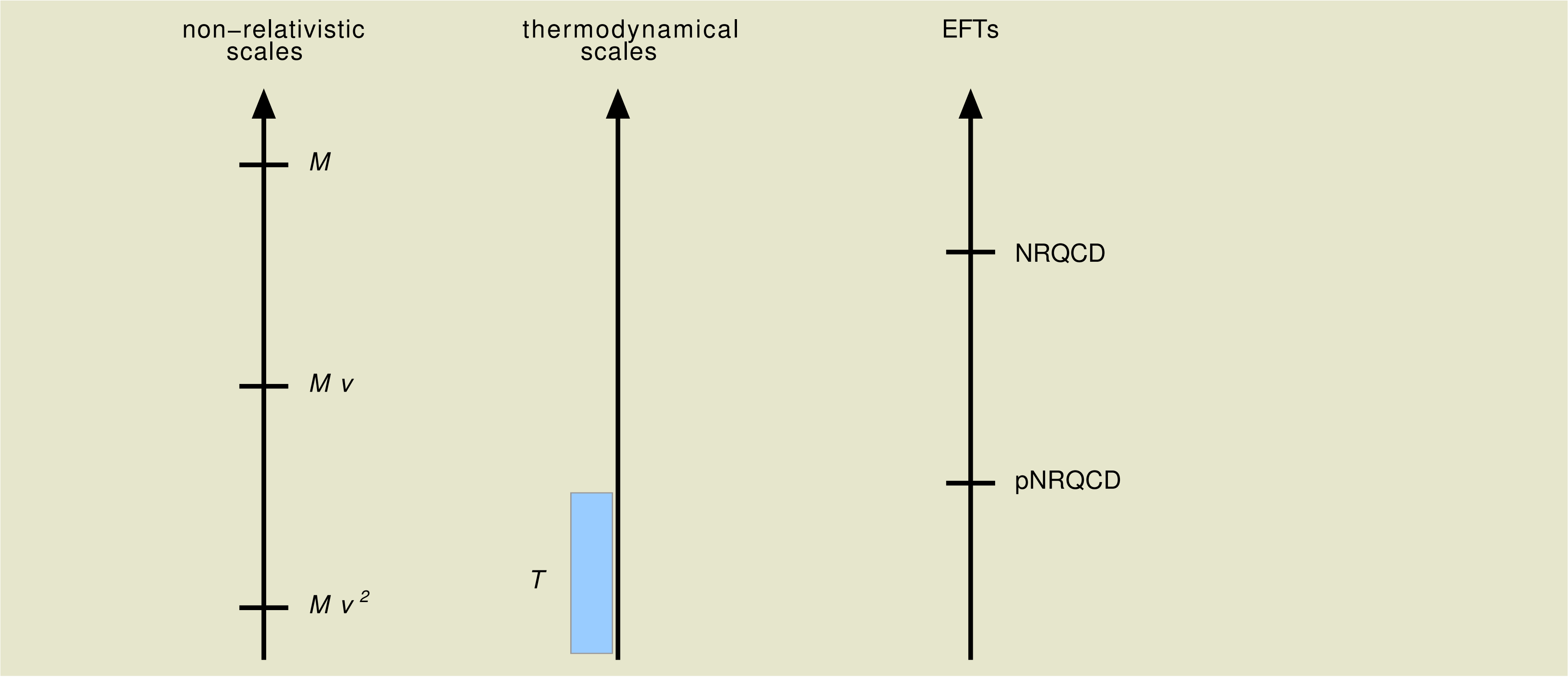
- the thermodynamical scales:
  T (temperature), ...

  T may stand for the inverse correlation length between system and environment. For definiteness we will assume that the system is locally in thermal equilibrium so that a slowly varying time-dependent temperature can be defined.

The non-relativistic scales are hierarchically ordered: $M \gg Mv \gg Mv^2$

Non-relativistic EFTs of QCD

The existence of a hierarchy of energy scales calls for a description of the system in terms of a hierarchy of EFTs.



- Brambilla Pineda Soto Vairo RMP 77 (2005) 1423
  Brambilla Ghiglieri Petreczky Vairo PRD 78 (2008) 014017

pNRQCD

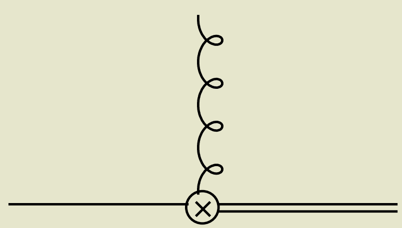
$$\mathcal{L} = -\frac{1}{4} F^a_{\mu\nu} F^{\mu\nu\, a} + \sum_{i=1}^{n_f} \bar{q}_i \, i \not{D} q_i + \int d^3r \mathrm{Tr} \left\{ \mathrm{S}^\dagger \left(i\partial_0 - h_s \right) \mathrm{S} + \mathrm{O}^\dagger \left(iD_0 - h_o \right) \mathrm{O} \right\}$$

- LO in r

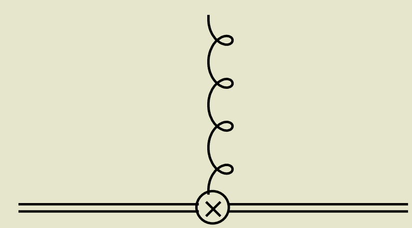
$\theta(T)\, e^{-iTh_s}$ $\qquad$ $\theta(T)\, e^{-iTh_o} \left(e^{-i \int dt\, A^{\mathrm{adj}}} \right)$

$$+V_A \mathrm{Tr} \left\{ \mathrm{O}^\dagger \mathbf{r} \cdot g\mathbf{E}\, \mathrm{S} + \mathrm{S}^\dagger \mathbf{r} \cdot g\mathbf{E}\, \mathrm{O} \right\} + \frac{V_B}{2} \mathrm{Tr} \left\{ \mathrm{O}^\dagger \mathbf{r} \cdot g\mathbf{E}\, \mathrm{O} + \mathrm{O}^\dagger \mathrm{O} \mathbf{r} \cdot g\mathbf{E} \right\}$$

- NLO in r



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



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

$$+ \cdots$$

Quarkonium in a fireball

- After the heavy-ion collisions, heavy quark-antiquarks propagate freely up to 0.6 fm.
- From 0.6 fm to the freeze-out time t_F they propagate in the medium.
- We assume the medium infinite, homogeneous and isotropic.
- We assume the heavy quarks comoving with the medium.
- We assume the medium to be locally in thermal equilibrium,

  i.e., the temperature T of the medium changes (slowly) with time:

$$T = T_0 \left(\frac{t_0}{t}\right)^{v_s^2}, \quad t_0 = 0.6 \text{ fm}, \quad v_s^2 = \frac{1}{3} \text{ (sound velocity)}$$

◦ Bjorken PRD 27 (1983) 140

Fireball's initial temperature

The initial temperature T_0 may account for different centralities

centrality (%)	$\langle b \rangle$ (fm)	T_0 (MeV)
$0 - 20$	4.76	466
$0 - 10$	3.4	471
$10 - 20$	6.0	461
$20 - 90$	11.6	360
$20 - 30$	7.8	449
$30 - 40$	9.35	433
$40 - 50$	10.6	412
$30 - 50$	9.9	425
$50 - 70$	12.2	366
$50 - 100$	13.6	304

- CMS PRC 84 (2011) 024906
  Miller Reygers Sanders Steinberg ARNPS 57 (2007) 205

Quarkonium as a Coulombic bound state

The lowest quarkonium states (1S bottomonium and charmonium, 2S bottomonium) are the most tightly bound. For these we assume the hierarchy of energy scales

$$M \gg \frac{1}{r} \sim M\alpha_{\rm s} \gg T \sim gT \gg \text{any other scale}, \qquad\qquad v \sim \alpha_{\rm s}$$

This qualifies the bound state as Coulombic:

- quark-antiquark color singlet Hamiltonian $= h_s = \dfrac{\mathbf{p}^2}{M} - \dfrac{4}{3}\dfrac{\alpha_{\rm s}}{r}$
- quark-antiquark color octet Hamiltonian $= h_o = \dfrac{\mathbf{p}^2}{M} + \dfrac{\alpha_{\rm s}}{6r}$

The octet potential describes an unbound quark-antiquark pair.

Density matrices

- Subsystem: heavy quarks/quarkonium
- Environment: quark gluon plasma

We may define a density matrix in pNRQCD for the heavy quark-antiquark pair in a singlet and octet configuration:

$$
\begin{aligned}
\langle \mathbf{r}', \mathbf{R}' | \rho_s(t';t) | \mathbf{r}, \mathbf{R} \rangle &\equiv \mathrm{Tr}\{\rho_{\mathrm{full}}(t_0)\, S^\dagger(t, \mathbf{r}, \mathbf{R}) S(t', \mathbf{r}', \mathbf{R}')\} \\
\langle \mathbf{r}', \mathbf{R}' | \rho_o(t';t) | \mathbf{r}, \mathbf{R} \rangle \frac{\delta^{ab}}{8} &\equiv \mathrm{Tr}\{\rho_{\mathrm{full}}(t_0)\, O^{a\,\dagger}(t, \mathbf{r}, \mathbf{R}) O^b(t', \mathbf{r}', \mathbf{R}')\}
\end{aligned}
$$

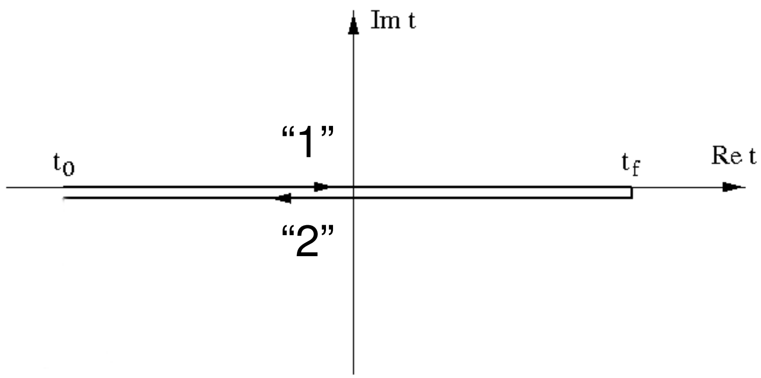
$t_0 \approx 0.6$ fm is the time formation of the plasma.

The system is in non-equilibrium because through interaction with the environment (quark gluon plasma) singlet and octet quark-antiquark states continuously transform in each other although the number of heavy quarks is conserved: $\mathrm{Tr}\{\rho_s\} + \mathrm{Tr}\{\rho_o\} = 1$.

Closed-time path formalism

In the closed-time path formalism we can represent the density matrices as 12 propagators on a closed time path:

$$
\begin{aligned}
\langle \mathbf{r}', \mathbf{R}' | \rho_s(t';t) | \mathbf{r}, \mathbf{R} \rangle &= \langle S_1(t', \mathbf{r}', \mathbf{R}') S_2^\dagger(t, \mathbf{r}, \mathbf{R}) \rangle \\
\langle \mathbf{r}', \mathbf{R}' | \rho_o(t';t) | \mathbf{r}, \mathbf{R} \rangle \frac{\delta^{ab}}{8} &= \langle O_1^b(t', \mathbf{r}', \mathbf{R}') O_2^{a\dagger}(t, \mathbf{r}, \mathbf{R}) \rangle
\end{aligned}
$$



Differently from the thermal equilibrium case 12 propagators are relevant (in thermal equilibrium they are exponentially suppressed).

12 propagators are not time ordered, while 11 and 22 operators select the forward time direction $\propto \theta(t-t')$, $\theta(t'-t)$.

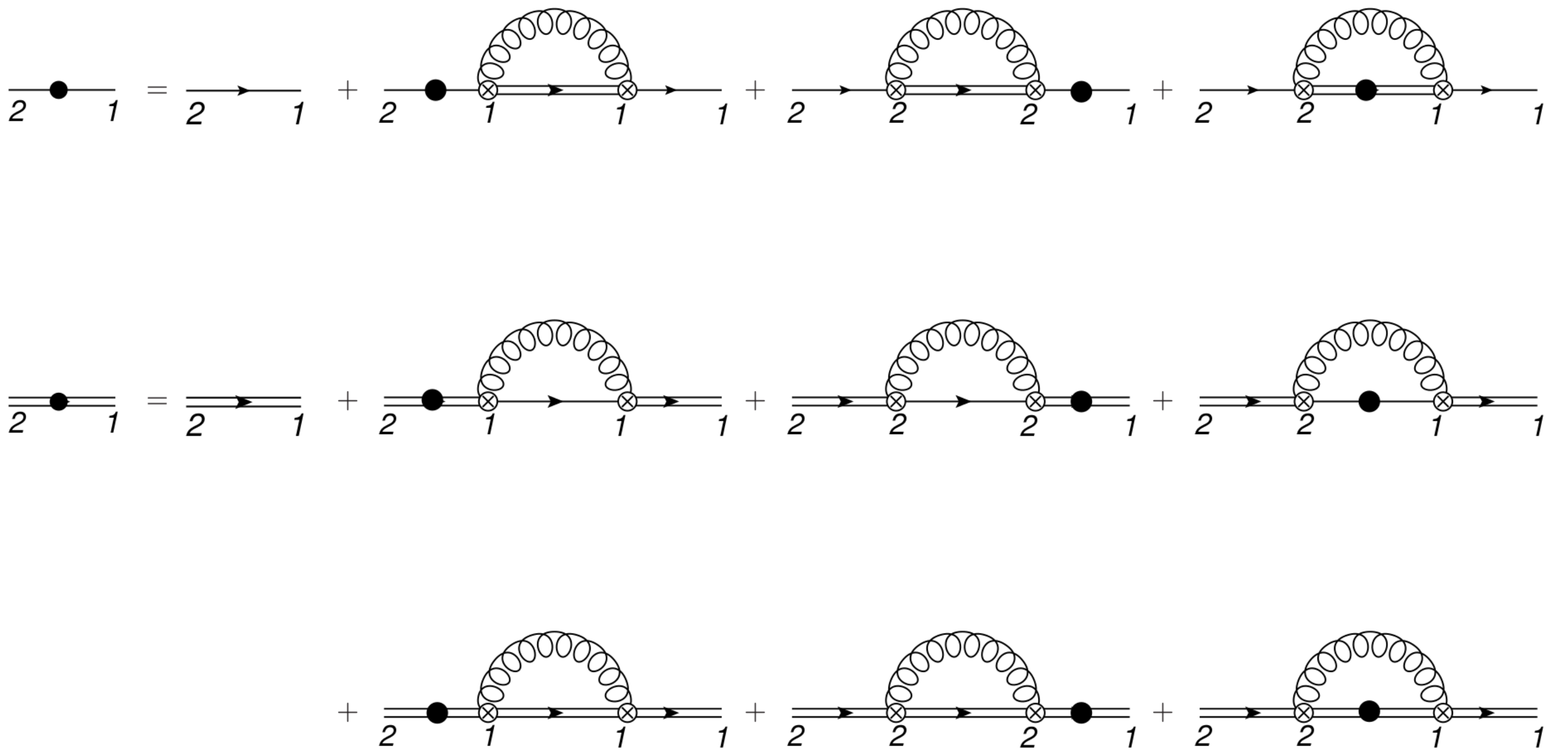
Expansions

- The density of heavy quarks is much smaller than the one of light quarks: we expand at first order in the heavy quark-antiquark density.
- We consider T much smaller than the Bohr radius of the quarkonium: we expand up to order r^2 in the multipole expansion.

The evolution depends on the density at initial time: non Markovian evolution.

Resummation

Resumming $(t - t_0)\times$ self-energy contributions à la Schwinger–Dyson ...



The resummation is accurate at order r^2 and consistent with unitary evolution at leading order.

Evolution equations I

... and differentiating over time we obtain the coupled evolution equations:

$$
\begin{aligned}
\frac{d\rho_s(t;t)}{dt} &= -i[h_s, \rho_s(t;t)] - \Sigma_s(t)\rho_s(t;t) - \rho_s(t;t)\Sigma_s^\dagger(t) + \Xi_{so}(\rho_o(t;t), t) \\
\frac{d\rho_o(t;t)}{dt} &= -i[h_o, \rho_o(t;t)] - \Sigma_o(t)\rho_o(t;t) - \rho_o(t;t)\Sigma_o^\dagger(t) + \Xi_{os}(\rho_s(t;t), t) \\
&\quad + \Xi_{oo}(\rho_o(t;t), t)
\end{aligned}
$$

- The evolution equations are now valid for large time.
- The evolution equations are Markovian.

Interpretation

- The self energies Σ_s and Σ_o provide the in-medium induced mass shifts, $\delta m_{s,o}$, and widths, $\Gamma_{s,o}$, for the color-singlet and color-octet heavy quark-antiquark systems respectively:

$$
\begin{aligned}
-i\Sigma_{s,o}(t) + i\Sigma^{\dagger}_{s,o}(t) &= 2\,\mathrm{Re}\,(-i\Sigma_{s,o}(t)) = 2\delta m_{s,o}(t) \\
\Sigma_{s,o}(t) + \Sigma^{\dagger}_{s,o}(t) &= -2\,\mathrm{Im}\,(-i\Sigma_{s,o}(t)) = \Gamma_{s,o}(t)
\end{aligned}
$$

- Ξ_{so} accounts for the production of singlets through the decay of octets, and Ξ_{os} and Ξ_{oo} account for the production of octets through the decays of singlets and octets respectively. There are two octet production mechanisms/octet chromoelectric dipole vertices in the pNRQCD Lagrangian.
- The conservation of the trace of the sum of the densities, i.e., the conservation of the number of heavy quarks, follows from

$$
\begin{aligned}
\mathrm{Tr}\left\{\rho_s(t;t)\left(\Sigma_s(t) + \Sigma^{\dagger}_s(t)\right)\right\} &= \mathrm{Tr}\left\{\Xi_{os}(\rho_s(t;t),t)\right\} \\
\mathrm{Tr}\left\{\rho_o(t;t)\left(\Sigma_o(t) + \Sigma^{\dagger}_o(t)\right)\right\} &= \mathrm{Tr}\left\{\Xi_{so}(\rho_o(t;t),t) + \Xi_{oo}(\rho_o(t;t),t)\right\}
\end{aligned}
$$

Evolution equations II

An alternative way of writing the evolution equations is

$$\frac{d\rho}{dt} = -i[H,\rho] + \sum_{nm} h_{nm} \left(L_i^n \rho L_i^{m\,\dagger} - \frac{1}{2}\{L_i^{m\,\dagger} L_i^n, \rho\} \right)$$

$$\rho = \begin{pmatrix} \rho_s & 0 \\ 0 & \rho_o \end{pmatrix} \qquad H = \begin{pmatrix} h_s + \frac{\Sigma_s - \Sigma_s^\dagger}{2i} & 0 \\ 0 & h_o + \frac{\Sigma_o - \Sigma_o^\dagger}{2i} \end{pmatrix}$$

$$\Sigma_s(t) = r^i\, A_i^{so\,\dagger}(t) \qquad \Sigma_o(t) = \frac{r^i\, A_i^{os\,\dagger}(t)}{8} + \frac{5}{16} r^i\, A_i^{oo\,\dagger}(t)$$

$$L_i^0 = \begin{pmatrix} 0 & 0 \\ 0 & 1 \end{pmatrix} r^i \qquad L_i^1 = \begin{pmatrix} 0 & 0 \\ 0 & \frac{5}{16} A_i^{oo\,\dagger} \end{pmatrix}$$

$$L_i^2 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix} r^i \qquad L_i^3 = \begin{pmatrix} 0 & \frac{1}{8} A_i^{os\,\dagger} \\ A_i^{so\,\dagger} & 0 \end{pmatrix}$$

with $A_i^{so}(t) = \frac{g^2}{6} \int_{t_0}^{t} dt_2\, e^{ih_s(t_2 - t)}\, r^j\, e^{ih_o(t-t_2)}\, \langle E^{a,j}(t_2,\mathbf{0}) E^{a,i}(t,\mathbf{0}) \rangle$

Positivity

The matrix h_{nm} is

$$h = \begin{pmatrix} 0 & 1 & 0 & 0 \\ 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & 1 & 0 \end{pmatrix} .$$

If h were a positive definite matrix then it would always be possible to redefine the operators L_i^n in such a way that the evolution equation would be of the Lindblad form.

Since, however, h is not a positive definite matrix, the Lindblad theorem does not guarantee that the equations may be brought into a Lindblad form.

From the evolution equations to the Lindblad equation

- We observe the evolution for large times ($\propto$ freeze-out time). Hence for $t - t_0$ larger than any other time scale that appears in the problem

$$\int_{t_0}^{t} dt_2\, f(t_2) \approx \int_0^{\infty} ds\, f(t-s)$$

- We assume the evolution to be quasistatic: $1/T \times dT/dt \sim 1/t \ll E$. The time dependence of the temperature can be neglected at leading order

$$\langle E^{a,i}(t,\mathbf{0}) E^{a,j}(t-s,\mathbf{0})\rangle \approx \langle E^{a,i}(s,\mathbf{0}) E^{a,j}(0,\mathbf{0})\rangle \approx \langle E^{a,i}(0,\mathbf{0}) E^{a,j}(-s,\mathbf{0})\rangle$$

An exact solution: the static limit case $1/r \gg T \gg E$

We consider the evolution equations for static quarks under the condition

$$\frac{1}{r} \gg T \gg E \sim \frac{\alpha_{\rm s}}{r}$$

Since $T \gg E$, energy-dependent exponentials can be set equal to one:

$$\begin{aligned}
(\Sigma_s + \Sigma_s^\dagger)\rho_s &= \Xi_{os} \equiv \Gamma_s \rho_s \\
(\Sigma_o + \Sigma_o^\dagger)\rho_o - \Xi_{oo} &= \Xi_{so} = \frac{1}{8}(\Sigma_s + \Sigma_s^\dagger)\rho_o \equiv \frac{\Gamma_s}{8}\rho_o
\end{aligned}$$

The evolution equations read

$$\begin{aligned}
\frac{d\rho_{s\,r}(t;t)}{dt} &= \Gamma_s(t)\left[\frac{\rho_{o\,r}(t;t)}{8} - \rho_{s\,r}(t;t)\right] \\
\frac{d\rho_{o\,r}(t;t)}{dt} &= -\Gamma_s(t)\left[\frac{\rho_{o\,r}(t;t)}{8} - \rho_{s\,r}(t;t)\right]
\end{aligned}$$

An exact solution: the static limit case $1/r \gg T \gg E$

The initial condition describes quark-antiquark pair at distance $\mathbf{r}$ with arbitrary color:

$$\rho_s(t_0;t_0) = \rho_{s\,r}(t_0;t_0)|\mathbf{r}\rangle\langle\mathbf{r}| \qquad \rho_o(t_0;t_0) = \rho_{o\,r}(t_0;t_0)|\mathbf{r}\rangle\langle\mathbf{r}|$$

The solution is

$$\begin{aligned}
\rho_{s\,r}(t;t) &= \frac{\rho_{s\,r}(t_0;t_0)}{9}\left[1+8e^{-\int_{t_0}^{t}\frac{dt'}{u(t')}}\right] + \frac{\rho_{o\,r}(t_0;t_0)}{9}\left[1-e^{-\int_{t_0}^{t}\frac{dt'}{u(t')}}\right] \\
\frac{\rho_{o\,r}(t;t)}{8} &= \frac{\rho_{s\,r}(t_0;t_0)}{9}\left[1-e^{-\int_{t_0}^{t}\frac{dt'}{u(t')}}\right] + \frac{\rho_{o\,r}(t_0;t_0)}{9}\left[1+\frac{e^{-\int_{t_0}^{t}\frac{dt'}{u(t')}}}{8}\right]
\end{aligned}$$

with $u(t) = \dfrac{8}{9}\dfrac{1}{\Gamma_s(t)}$.

- $t - t_0 \ll u(t)$ the thermal medium has a small impact on the distribution of quarkonia.
- $t - t_0 \gg u(t)$ the density approaches the large-time asymptotic value.

The static limit case without equilibration

We consider the general case

$$\Gamma_s(T) = \Gamma_s(T_0) \left(\frac{T}{T_0} \right)^n$$

If $nv_s^2 > 1$ the color-singlet density never reaches the thermal equilibrium value $1/9$. Instead it approaches the value

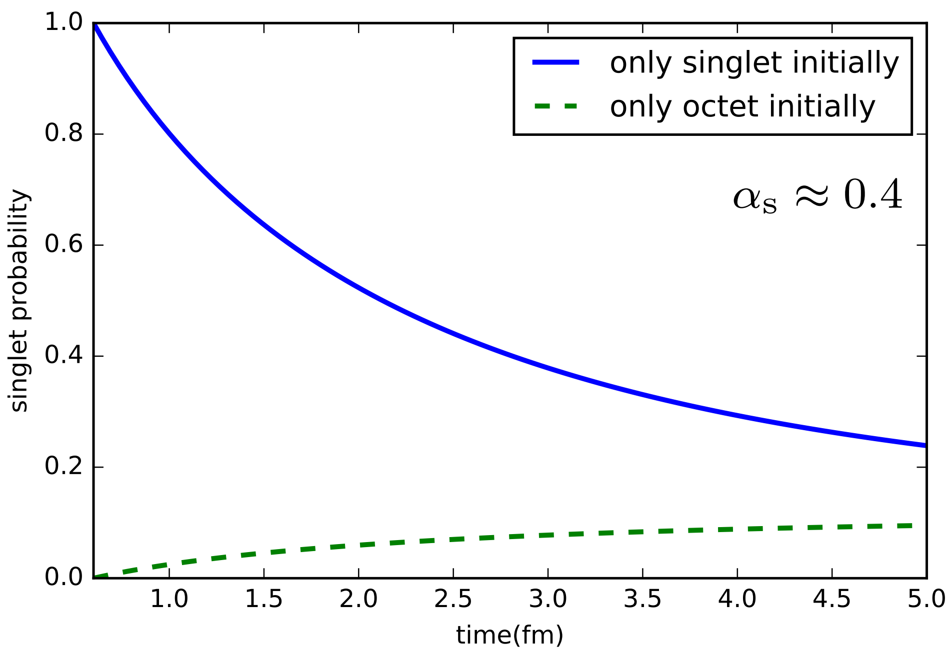
$$\frac{1}{9} - \frac{8}{9} \left(\frac{\rho_{o\,r}(t_0; t_0)}{8} - \rho_{s\,r}(t_0; t_0) \right) e^{\frac{9}{8} \, \frac{\Gamma_s(T_0) t_0}{1 - n v_s^2}}$$

The decrease with time of the decay width in the fireball is so fast that the static quark-antiquark densities do not have time to equilibrate.

The static limit case with equilibration

If $nv_s^2 < 1$, the color-singlet and color-octet densities reach their thermal equilibrium values exponentially fast. This is the situation realized by static quarks and antiquarks in the weakly-coupled plasma in the case $\alpha_{\rm s}/r \gg m_D$, for which $n = 1$:

$$\Gamma_s(t) = \Sigma_s(t) + \Sigma_s^\dagger(t) = 4\,\alpha_{\rm s}^3\,T$$



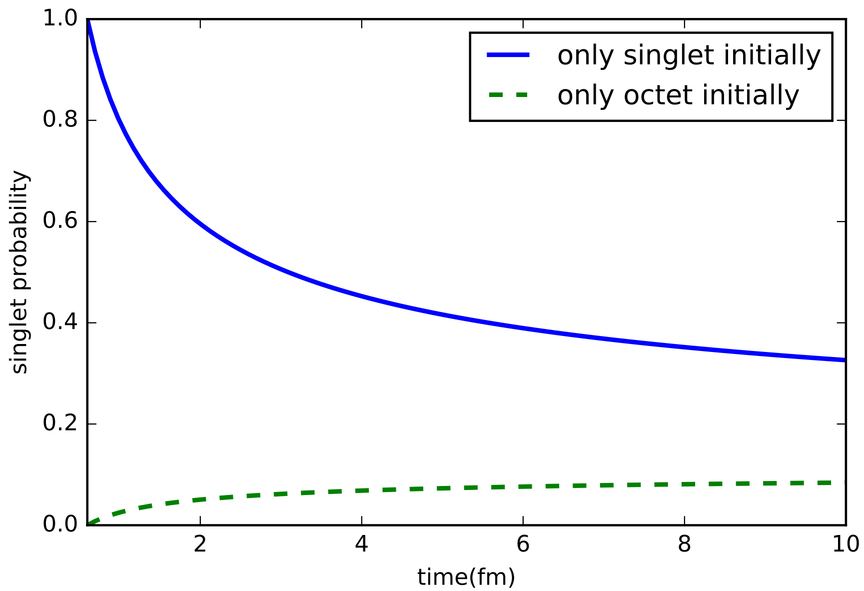
The static limit case with power-like equilibration

If $nv_s^2 = 1$ we have

$$\rho_{s\,r}(t;t) = 1 - \rho_{o\,r}(t;t) = \frac{1}{9} - \frac{8}{9}\left(\frac{\rho_{o\,r}(t_0;t_0)}{8} - \rho_{s\,r}(t_0;t_0)\right)\left(\frac{t_0}{t}\right)^{\frac{9}{8}\Gamma_s(T_0)t_0}$$

Color-singlet and color-octet densities reach their thermal equilibrium values with a power-like falloff. This situation is realized by static quarks and antiquarks in the strongly-coupled plasma, $T \sim m_D$, for which $n = 3$:

$$\Gamma_s(t) = \Sigma_s(t) + \Sigma_s^\dagger(t) = \kappa(t)\, r^2 \sim T^3 r^2$$



Results:

The Out of Equilibrium Evolution

Heavy quark-antiquarks in a strongly coupled medium: $T \gg E$

If $E \ll T \sim m_D$ the Lindblad equation for a strongly coupled plasma reads

$$\frac{d\rho}{dt} = -i[H,\rho] + \sum_i \left(C_i \rho C_i^\dagger - \frac{1}{2}\{C_i^\dagger C_i, \rho\}\right)$$

$$\rho = \begin{pmatrix} \rho_s & 0 \\ 0 & \rho_o \end{pmatrix}$$

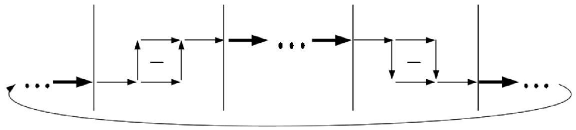
$$H = \begin{pmatrix} h_s & 0 \\ 0 & h_o \end{pmatrix} + \frac{r^2}{2}\gamma(t) \begin{pmatrix} 1 & 0 \\ 0 & \frac{7}{16} \end{pmatrix},$$

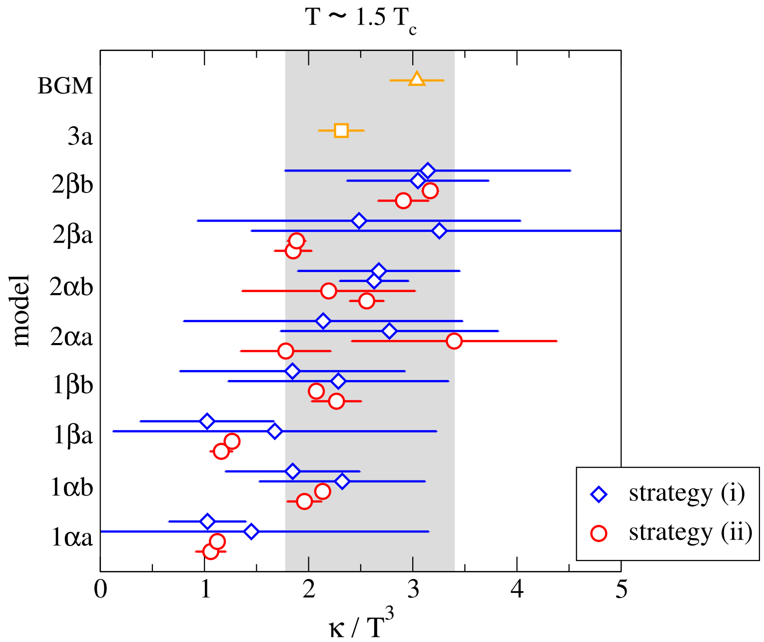
$$C_i^0 = \sqrt{\frac{\kappa(t)}{8}}\, r^i \begin{pmatrix} 0 & 1 \\ \sqrt{8} & 0 \end{pmatrix}, \qquad C_i^1 = \sqrt{\frac{5\kappa(t)}{16}}\, r^i \begin{pmatrix} 0 & 0 \\ 0 & 1 \end{pmatrix}$$

κ

Low energy parameters may be determined by numerical calculations in lattice QCD.
κ is the heavy-quark momentum diffusion coefficient:

$$\kappa = \frac{g^2}{18} \,\mathrm{Re} \int_{-\infty}^{+\infty} ds \,\langle\, \mathrm{T}\, E^{a,i}(s,\mathbf{0})\, \phi^{ab}(s,0)\, E^{b,i}(0,\mathbf{0}) \rangle =$$





○ Francis Kaczmarek Laine Neuhaus Ohno PRD 92 (2015) 116003

γ

$$\gamma = \frac{g^2}{18}\,\mathrm{Im}\int_{-\infty}^{+\infty} ds\,\langle\,\mathrm{T}\,E^{a,i}(s,\mathbf{0})\,\phi^{ab}(s,0)\,E^{b,i}(0,\mathbf{0})\rangle$$

γ is known only in perturbation theory:

$$\gamma = -4\zeta(3)\frac{\alpha_{\mathrm{s}}}{\pi}Tm_D^2 + \frac{16}{3}\zeta(3)\alpha_{\mathrm{s}}^2T^3$$

A value that at leading order is negative.

- Brambilla Ghiglieri Petreczky Vairo PRD 78 (2008) 014017

Initial conditions

- The production of singlets is α_s suppressed compared to that of octets.
  - `Cho Leibovich PRD 53 (1996) 6203`

- Our choice at $t = 0$ is

$$\rho_s = N|\mathbf{0}\rangle\langle\mathbf{0}|, \qquad \rho_o = \frac{\delta}{\alpha_s(M)}\rho_s$$

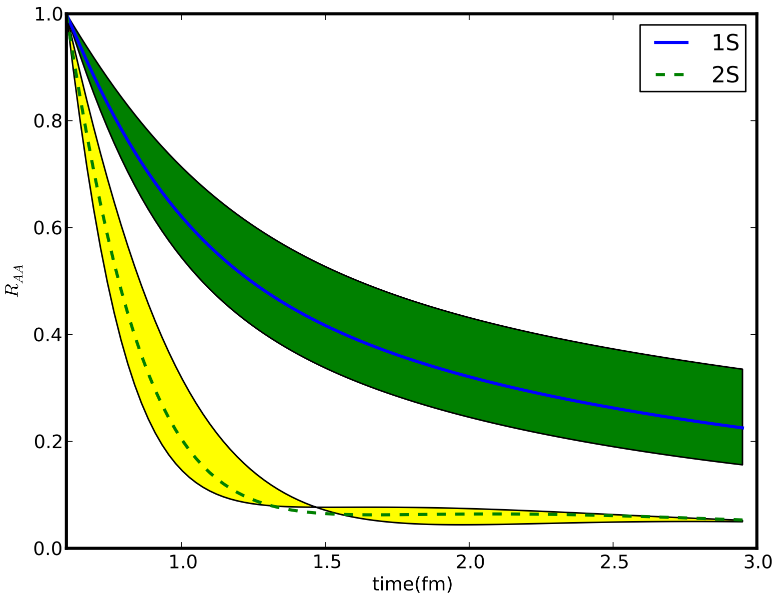
N is fixed by $\mathrm{Tr}\{\rho_s\} + \mathrm{Tr}\{\rho_o\} = 1$
δ fixes the octet fraction with respect to the singlet.

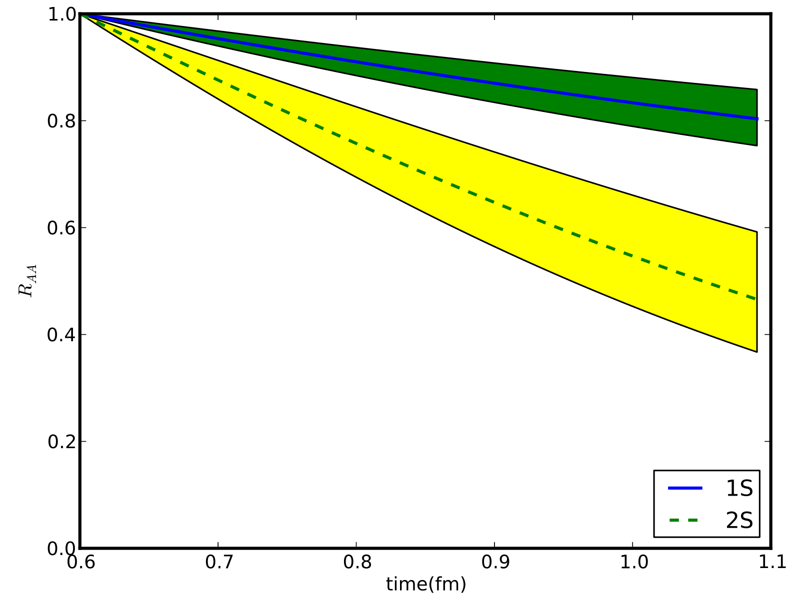
Nuclear modification factor

We compute the nuclear modification factor R_{AA}:

$$R_{AA}(nS) = \frac{\langle n, \mathbf{q}|\rho_s(t_F; t_F)|n, \mathbf{q}\rangle}{\langle n, \mathbf{q}|\rho_s(0; 0)|n, \mathbf{q}\rangle}$$

Time evolution of R_{AA} for $1S$ and $2S$ bottomonium

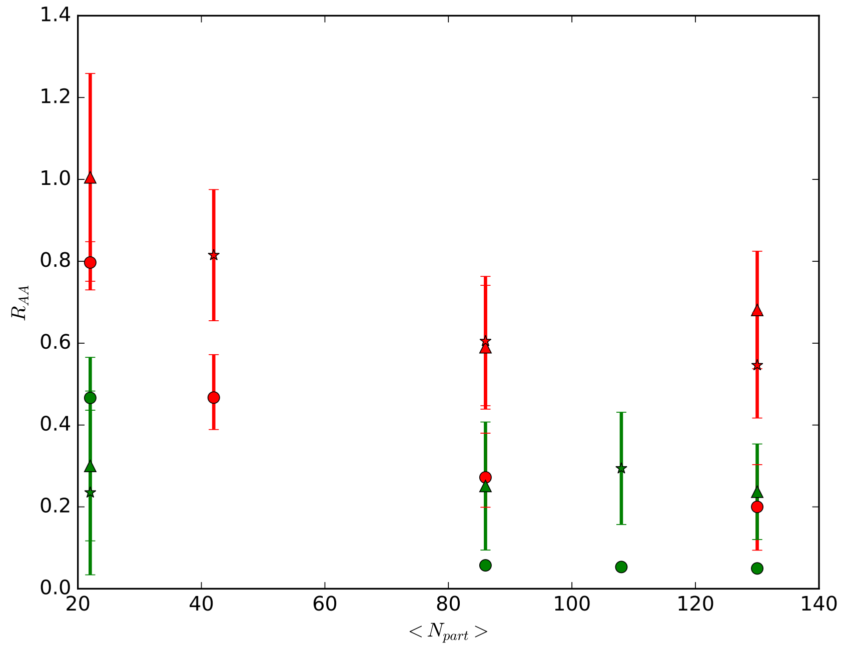




$30 - 50\%$ centrality (left) and $50 - 100\%$ centrality (right)

$\gamma = 0$ and $\delta = 1$

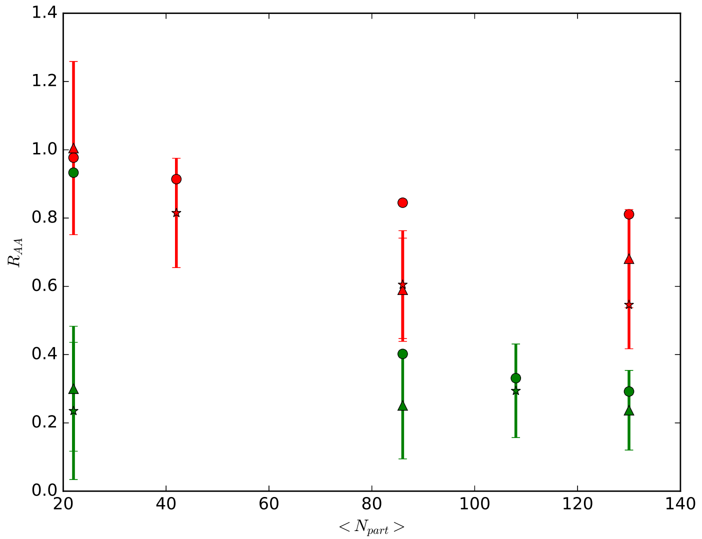
Bottomonium nuclear modification factor vs CMS data



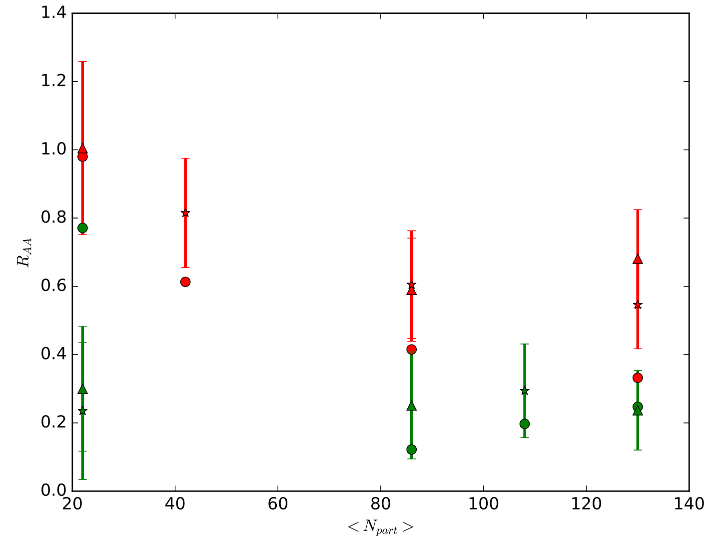
○: theory results; △: CMS data; red: $\Upsilon(1S)$; green: $\Upsilon(2S)$.

- data from CMS coll. PRL 109 (2012) 222301 and PLB 770 (2017) 357

Bottomonium nuclear modification factor vs CMS data



$\kappa/T^3 = 0.25$ and $\gamma = 0$;



$\kappa/T^3 = 2.6$ and $\gamma/T^3 = 6$.

Conclusions

We have shown how the heavy quark-antiquark pair out-of-equilibrium evolution can be treated in the framework of pNRQCD. With respect to previous determinations:

- the medium may be a strongly-coupled plasma (not necessarily a quark-gluon plasma) whose characteristics are determined by lattice calculations;
- the total number of heavy quarks, i.e., $\mathrm{Tr}\{\rho_s\} + \mathrm{Tr}\{\rho_o\}$, is preserved by the evolution equations;
- the non-abelian nature of QCD is fully accounted for;
- the treatment does not rely on classical approximations.

The evolution equations follow from assuming the inverse size of the quark-antiquark system to be larger than any other scale of the medium and from being accurate at first non-trivial order in the multipole expansion and at first order in the heavy-quark density.

Under some conditions (large time, quasistatic evolution, temperature much larger than the inverse evolution time of the quarkonium) the evolution equations are of the Lindblad form. Their numerical solution provides $R_{AA}(nS)$ close to experimental data.

Outlook

Several improvements are possible starting from this first analysis, e.g.,

- use more refined models for the hydrodynamical evolution than the Bjorken one;
- use more refined initial conditions accounting for higher-order production mechanisms, or for state of matters, like the color glass condensate;
- include the momentum dependence of the quark-antiquark pairs;
- determine κ and γ on unquenched lattices for wide ranges of temperature;
- relax the conditions leading to the Lindblad form and solve directly the non-Lindblad evolution equations;
- ...