



# Heavy quarkonium in a moving thermal bath

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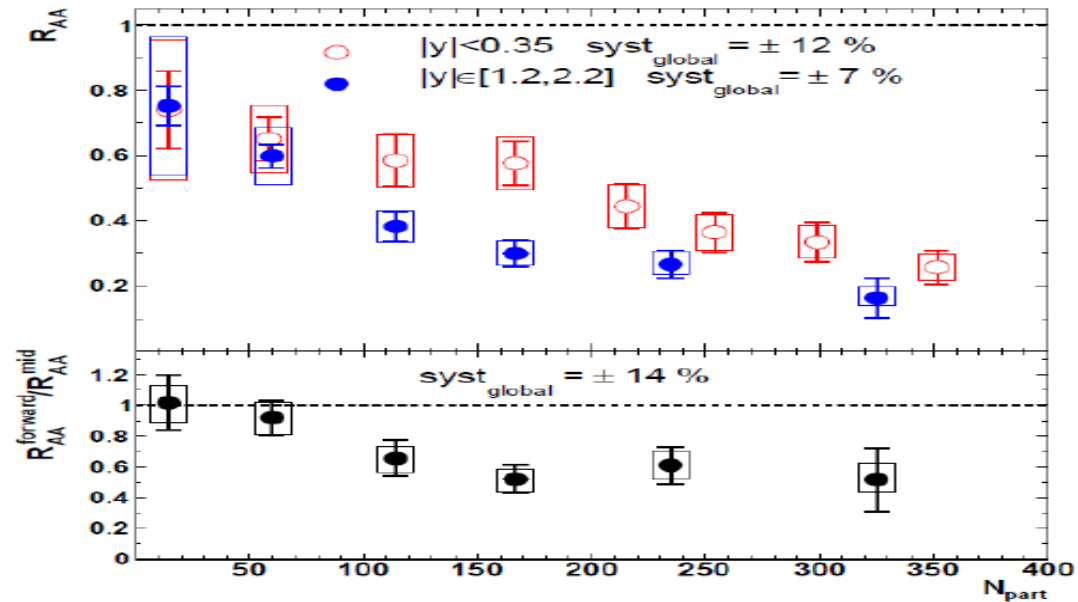
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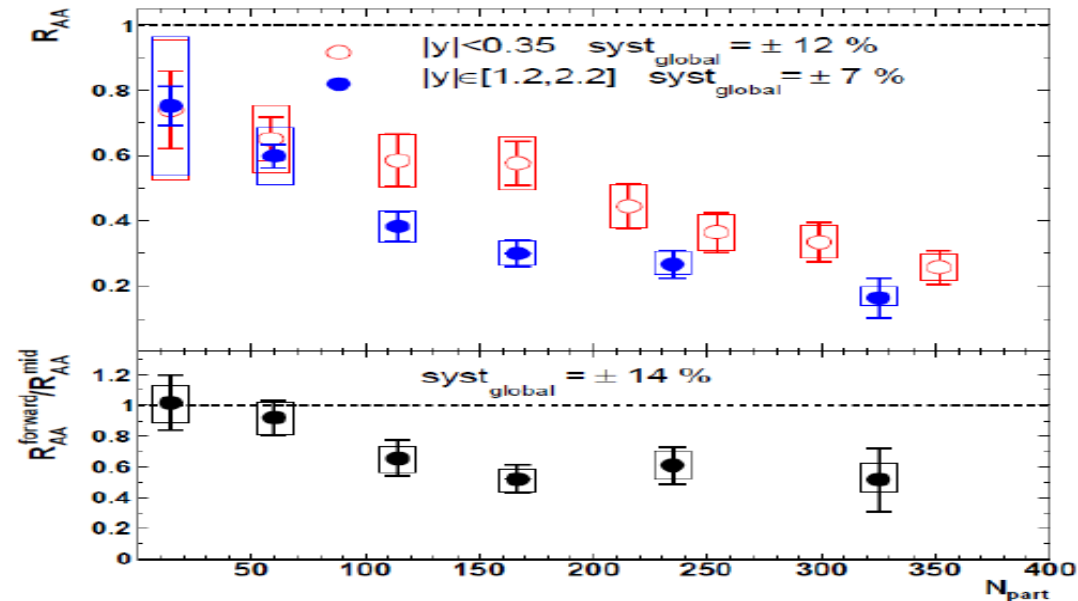
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    - Heavy quarkonium at weak coupling ( $\Upsilon(1S)$ ,  $J/\psi$ )
    - Weakly coupled QGP
  - Fruitful strategy in the case at rest (Miguel Angel Escobedo, JS; Phys. Rev. A 78, 032520 (2008), Phys. Rev. A 82, 042506 (2010))

# Motivation



(H. Pereira Da Costa, for the PHENIX collaboration, arXiv:1007.3688)

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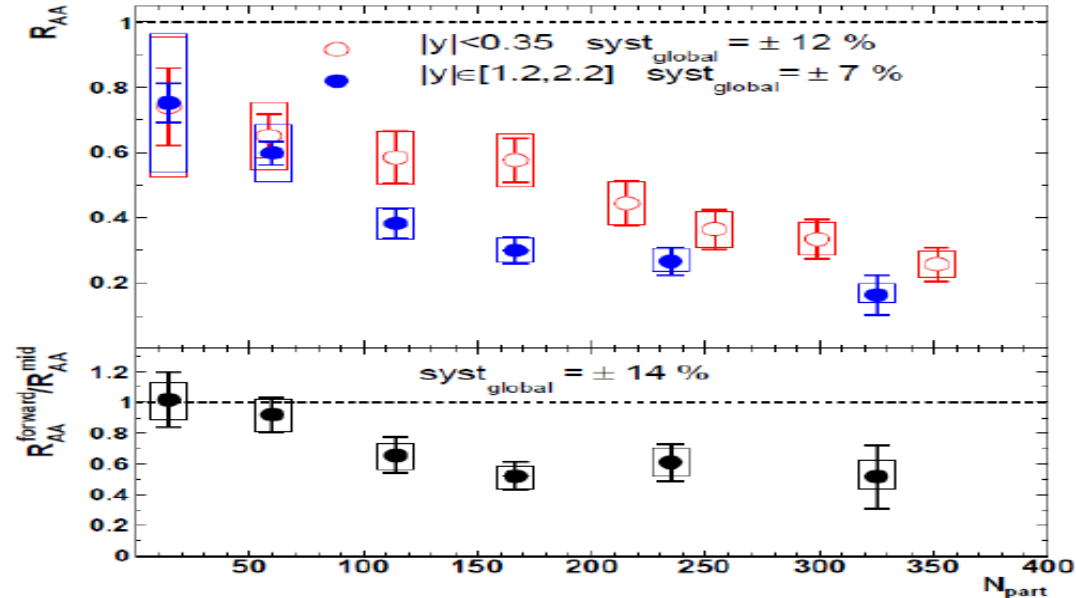


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# Motivation



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- The  $J/\psi$  suppression depends on the rapidity
- Is this due, at least in part, to the fact that the in-medium properties of  $J/\psi$  depend on the velocity?

# Relevant scales (hydrogen atom)



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- $m \neq 0, T = 0$  case:
  - $m$  (hard), electron mass
  - $m\alpha/n$  (soft), inverse Bohr radius,  $\alpha = e^2/4\pi; e$ , electron charge
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- $m \neq 0, T \neq 0$  case: what is the interplay among the scales above?

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•  $m \neq 0, T \neq 0$  case: contributions of energies above  $T$  are exponentially suppressed by Boltzmann factors

# Non-Relativistic QED (T=0)

$$\begin{aligned}\mathcal{L}_{NRQED} = & -\frac{1}{4}d_1 F_{\mu\nu}F^{\mu\nu} + \frac{d_2}{m^2}F_{\mu\nu}D^2F^{\mu\nu} + N^\dagger iD^0N + \\ & +\psi^\dagger\left(iD^0 + \frac{\mathbf{D}^2}{2m} + \frac{\mathbf{D}^4}{8m^3} + c_F e \frac{\boldsymbol{\sigma}\mathbf{B}}{2m} + c_D e \frac{\nabla\mathbf{E}}{8m^2} + \right. \\ & \left. + ic_S e \frac{\boldsymbol{\sigma}(\mathbf{D}\times\mathbf{E} - \mathbf{E}\times\mathbf{D})}{8m^2}\right)\psi\end{aligned}$$

(Caswell, Lepage, 1986)

# Potential NRQED (T=0)



$$\begin{aligned} L_{pNRQED} = & - \int d^3\mathbf{x} \frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \int d^3\mathbf{x} S^\dagger(t, \mathbf{x}) \left( iD_0 + \frac{\nabla^2}{2m} + \frac{Z\alpha}{|\mathbf{x}|} + \right. \\ & \left. + \frac{\nabla^4}{8m^3} + \frac{Ze^2}{m^2} \left( -\frac{c_D}{8} + 4d_2 \right) \delta^3(\mathbf{x}) + ic_S \frac{Z\alpha}{4m^2} \boldsymbol{\sigma} \cdot \left( \frac{\mathbf{x}}{|\mathbf{x}|^3} \times \nabla \right) \right) S(t, \mathbf{x}) \\ & + \int d^3\mathbf{x} S^\dagger(t, \mathbf{x}) e\mathbf{x} \cdot \mathbf{E} S(t, \mathbf{x}). \end{aligned}$$

(Pineda, Soto, 1997)



# Hard Thermal Loops EFT ( $m=0$ )

$$\delta\mathcal{L}_{HTL} = \frac{1}{2}m_D^2 F^{\mu\alpha} \int \frac{d\Omega}{4\pi} \frac{k_\alpha k_\beta}{-(k.\partial)^2} F^{\mu\beta} + m_e^2 \bar{\psi} \gamma^\mu \int \frac{d\Omega}{4\pi} \frac{k_\mu}{k.\partial} \psi$$

$$k = (1, \hat{\mathbf{k}}), \quad m_D^2 = e^2 T^2 / 3, \quad m_e^2 = e^2 T^2 / 8$$

(Braaten, Pisarsky, 1992)

# The $v \neq 0$ case

- Bound state at rest, the medium moves at velocity  $v$  (Weldon, 82)

$$f(\beta k^0) \rightarrow f(\beta^\mu k_\mu) = \frac{1}{e^{|\beta^\mu k_\mu|} \pm 1}, \quad \beta^\mu = \frac{\gamma}{T}(1, \mathbf{v})$$

$$v = |\mathbf{v}|, \quad \gamma = 1/\sqrt{1 - v^2}$$

- $O(3)$  rotational symmetry is reduced to  $O(2)$
- In light cone coordinates  $k_+ = k_0 + k_3$ ,  $k_- = k_0 - k_3$

$$\beta^\mu k_\mu = \frac{1}{2} \left( \frac{k_+}{T_+} + \frac{k_-}{T_-} \right), \quad T_+ = T \sqrt{\frac{1+v}{1-v}}, \quad T_- = T \sqrt{\frac{1-v}{1+v}}$$

- For  $v \ll 1$  (moderate velocities),  $T_+ \sim T \sim T_-$
- For  $v \sim 1$  (relativistic velocities),  $T_+ \gg T \gg T_-$ 
  - Collinear region,  $k_+ \sim T_+$ ,  $k_- \sim T_-$
  - Soft (ultrasoft) region,  $k_+ \sim k_- \sim T_-$

# Moderate velocity ( $v \approx 1$ )

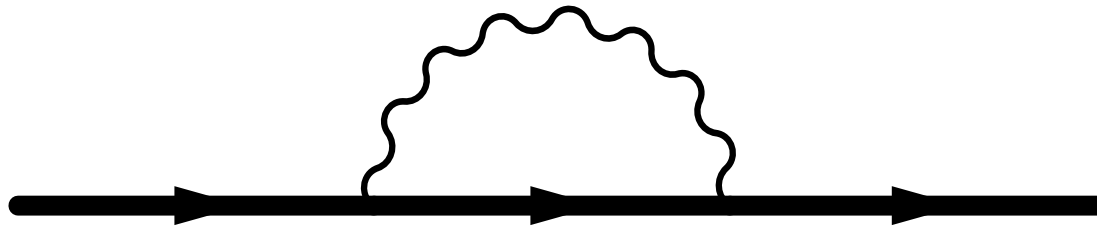
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  - The potentials remain the same as in the  $T = 0$  case
  - Thermal effects are encoded in the ultrasoft photons







$v \approx 1$ : the  $T \ll m\alpha/n$  case



## $v \approx 1$ : the $T \ll m\alpha/n$ case

• For  $T = \beta^{-1} \ll m\alpha^2/n^2$ , the result does not depend on  $v$ :

$$\begin{aligned}\delta E_n &= -\frac{4\pi^3\alpha}{45\beta^4} \langle n | \mathbf{x} \frac{\bar{P}_n}{(H_0 - E_n)} \mathbf{x} | n \rangle \left( 1 + \mathcal{O} \left( \left( \frac{n^2}{\beta m \alpha} \right)^2 \right) \right) \\ \delta \Gamma_n &= 0\end{aligned}$$

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$$\delta \Gamma_n = 0$$

- For  $T = \beta^{-1} \gg m\alpha^2/n^2$ ,  $l = 0$ :

$$\delta E_n = \frac{\alpha\pi T^2}{3m_e} - \frac{4Z\alpha^2}{3} \frac{|\phi_n(\mathbf{0})|^2}{m_e^2} \left( -\frac{1}{2v} \log \left( \frac{1+v}{1-v} \right) + 1 - \gamma \right)$$

$$+ \frac{2\alpha}{3\pi m_e^2} \sum_r |\langle n | \mathbf{p} | r \rangle|^2 (E_n - E_r) \log \left( \frac{2\pi T}{|E_n - E_r|} \right) \times \left( 1 + \mathcal{O} \left( \left( \frac{\beta m \alpha}{n^2} \right)^2 \right) \right)$$

$$\delta \Gamma_n = \frac{2Z^2\alpha^3 T \sqrt{1-v^2}}{3n^2 v} \log \left( \frac{1+v}{1-v} \right) \times \left( 1 + \mathcal{O} \left( \frac{\beta m \alpha}{n^2} \right) \right)$$

## $v \approx 1$ : the $T \ll m\alpha/n$ case

• For  $T = \beta^{-1} \gg m\alpha^2/n^2$ ,  $l \neq 0$ :

$$\delta E_{nlm} = \frac{\alpha\pi T^2}{3m_e} + \frac{2\alpha}{3\pi m_e^2} \sum_r |\langle n|\mathbf{p}|r\rangle|^2 (E_n - E_r) \log\left(\frac{-E_1}{|E_n - E_r|}\right) - \frac{Z^3\alpha^2 \langle 2l00|l0\rangle \langle 2l0m|lm\rangle}{2\pi m_e^2 a_0^3 n^3 l(l + \frac{1}{2})(l + 1)} \rho(v)$$

$$\delta\Gamma_{nlm} = \frac{Z^2\alpha^3 T \sqrt{1-v^2}}{3n^2 v} \left( 2 \log\left(\frac{1+v}{1-v}\right) + \left( \left(1 - \frac{3}{v^2}\right) \log\left(\frac{1+v}{1-v}\right) + \frac{6}{v} \right) \langle 2l00|l0\rangle \langle 2l0m|lm\rangle \right)$$

$$\rho(v) = \frac{1}{2v} \left(1 - \frac{1}{v^2}\right) \log\left(\frac{1+v}{1-v}\right) - \frac{2}{3} + \frac{1}{v^2}$$

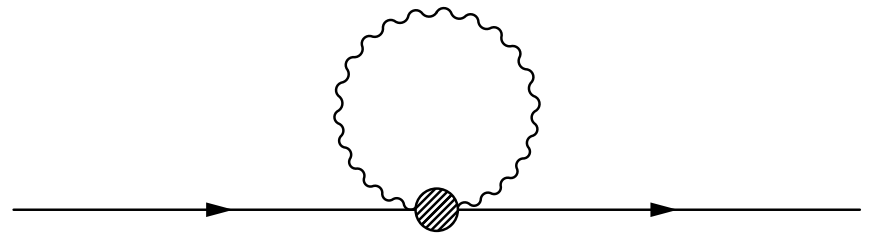
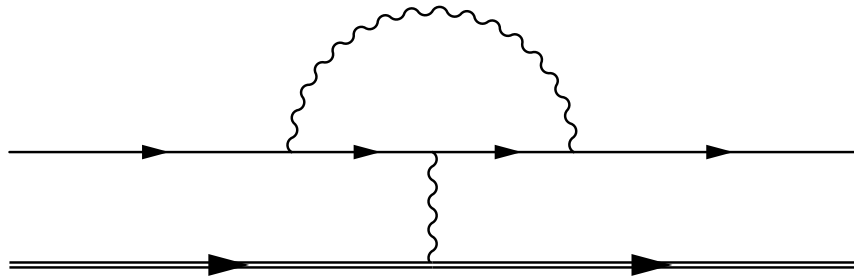
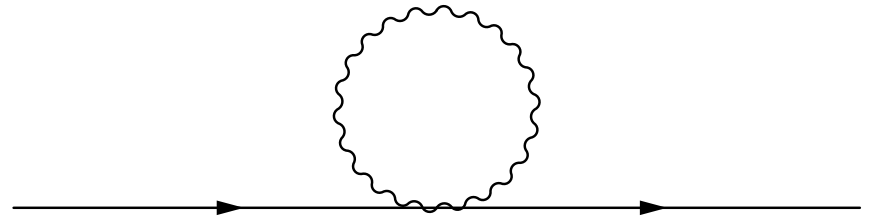
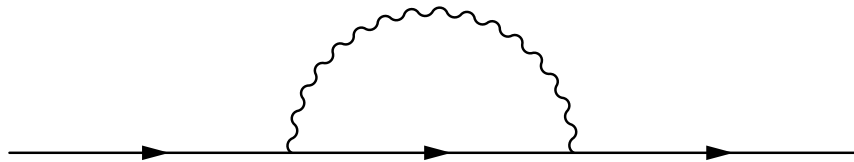
• The decay width decreases as  $v$  increases !

# Moderate velocity ( $v \approx 1$ )

- The  $T \ll m$  case: NRQED can be used as a starting point
  - The potentials depend on  $T$ :

$$\begin{aligned} \delta \mathcal{L}_{pNRQED} = & \int d^3 \mathbf{x} \left( \frac{\alpha \pi T^2}{3m_e} \psi^\dagger \psi - \frac{\pi \alpha T^2}{6m_e^3} \nabla \psi^\dagger \nabla \psi \right) \\ & + \int d^3 \mathbf{x}_1 d^3 \mathbf{x}_2 N^\dagger(t, \mathbf{x}_2) N(t, \mathbf{x}_2) \left[ -\frac{4Z\alpha}{3m_e^2} \left( \log \left( \frac{\mu}{2\pi T} \right) - \log 2 \right. \right. \\ & \left. \left. + \gamma + \frac{3}{8v} \left( 1 + \frac{1}{3v^2} \right) \log \left( \frac{1+v}{1-v} \right) - \frac{1}{4v^2} \right) \delta^3(\mathbf{x}_1 - \mathbf{x}_2) \right. \\ & \left. + \frac{\alpha \rho(v) v^i v^j \partial_{ij}^2 V_c(r)}{4\pi m_e^2 v^2} \right] \psi^\dagger(t, \mathbf{x}_1) \psi(t, \mathbf{x}_1) \end{aligned}$$

- $\mu$ , factorization scale arising from IR divergences, which cancels against the one of ultrasoft contributions in physical observables



## $v \approx 1$ : the $T \ll m$ case

- In the ultrasoft contributions,

$$1/(e^{|\beta^\mu k_\mu|} - 1) \rightarrow 1/|\beta^\mu k_\mu| - 1/2 + \dots$$

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$$\begin{aligned} \delta E_{nlm} = & \frac{\alpha\pi T^2}{3m_e} - \frac{\pi\alpha^3 T^2}{6m_e n^2} + \frac{4Z\alpha^2}{3m_e^2} \left( \gamma + \frac{3}{8v} \left( 1 + \frac{1}{3v^2} \right) - \frac{1}{4v^2} \right) |\phi_n(\mathbf{0})|^2 \\ & - \frac{\alpha\rho(v)v^i v^j}{4\pi m_e^2 v^2} \langle n | \partial_{ij}^2 V_c(r) | n \rangle + \frac{2\alpha}{3\pi m_e^2} \sum_r |\langle n | \mathbf{p} | r \rangle|^2 (E_n - E_r) \left( \log \left( \frac{2\pi T}{|E_n - E_r|} \right) + \frac{5}{6} \right) \end{aligned}$$

$$\begin{aligned} \delta \Gamma_{nlm} = & \frac{Z^2 \alpha^3 T \sqrt{1-v^2}}{3n^2 v} \left( 2 \log \left( \frac{1+v}{1-v} \right) \right. \\ & \left. + \left( \left( 1 - \frac{3}{v^2} \right) \log \left( \frac{1+v}{1-v} \right) + \frac{6}{v} \right) \langle 2l00 | l0 \rangle \langle 2l0m | lm \rangle \right) \end{aligned}$$

# Relativistic velocity ( $v \sim 1$ )

- The  $T_+ \sim m\alpha/n \gg T_- \gg m\alpha^2/n^2$  case: NRQED can be used as a starting point
- The collinear photons ( $k_+ \sim T_+$ ,  $k_- \sim T_-$ ) have virtualities  $\sim T^2 \ll (m\alpha/n)^2$  and hence must be kept in the effective theory: pNRQED  $\rightarrow$  pNRQED + SCET

$$\delta\mathcal{L}_{SCET} = c_1 \frac{\psi^\dagger \psi}{m_e} \frac{\bar{n}^\mu F_{\mu i}}{(\bar{n}\partial)} \frac{\bar{n}^\nu F_{\nu i}}{(\bar{n}\partial)} + c_2 \frac{\psi^\dagger \psi}{m_e} \frac{\bar{n}^\mu n^\nu F_{\mu\nu}}{(\bar{n}\partial)} \frac{\bar{n}^\alpha n^\beta F_{\alpha\beta}}{(\bar{n}\partial)} + \dots$$

$F_{\mu\nu}$ , collinear photons;  $\psi$ , NR electron field

$$n = (1, 0, 0, 1), \bar{n} = (1, 0, 0, -1)$$

- The pNRQED Lagrangian for ultrasoft photons remains the same

$v \sim 1$ : the  $T_+ \sim m\alpha/n \gg T_- \gg m\alpha^2/n^2$  case

• Calculation in pNRQED+ SCET

•  $\delta E_n^{\text{col}} = \frac{\pi\alpha T^2}{3m_e}, \quad \delta\Gamma_n^{\text{col}} = 0$

• There are three relevant regions in the us contribution:

•  $k_+, k_- \sim T_-$

•  $k_+, k_- \sim m\alpha^2/n^2$

•  $k_+ \sim m\alpha^2/n^2$  and  $k_- \sim m\alpha^2/n^2(T_-/T_+)$

$$\delta E_{n00}^{\text{us}} = -\frac{4Z\alpha^2}{3} \left( 1 - \gamma + \frac{1}{2} \log \left( \frac{1-v}{1+v} \right) \right) \frac{|\phi_n(0)|^2}{m_e^2}$$

$$- \frac{2\alpha}{3\pi m_e^2} \sum_r |\langle n|\mathbf{p}|r\rangle|^2 (E_n - E_r) \log \left( \frac{|E_n - E_r|}{2\pi T} \right)$$

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• Agreement with the  $v \rightarrow 1$  limit of the  $v \approx 1$  case

# Relativistic velocity ( $v \sim 1$ )

- The  $T_+ \sim m \gg m\alpha/n \gg T_- \gg m\alpha^2/n^2$  case: **QED** must be used as a starting point
  - The collinear photons ( $k_+ \sim T_+$ ,  $k_- \sim T_-$ ) have virtualities  $\sim T^2 \ll m^2$  and hence must be kept in the effective theory: **NRQED**  $\rightarrow$  **NRQED** + **SCET**
  - The collinear photons have virtualities  $\gg T_-^2$  and can be integrated out in the matching to **pNRQED**. They produce a global energy shift.
  - The **pNRQED** Lagrangian for ultrasoft photons remains the same
  - Agreement with the  $v \rightarrow 1$  limit of the  $v \approx 1$  case

# The $m\alpha/n \ll T \ll m$ case

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- One obtains a  $v$ -dependent **HTL** Lagrangian in the photon sector and  $v$ -dependent temperature corrections to the NRQCD matching coefficients
- One next matches to **pNRQED**, obtaining a  $v$  and  $T$  dependent potential



## The $m\alpha/n \ll T \ll m$ case at $v = 0$

$$V(r, T) = -\frac{\alpha e^{-m_D r}}{r} - \alpha m_D + i\alpha T \phi(m_D r) + \mathcal{O}\left(\frac{\alpha T^2}{m_\mu}\right)$$

$$\phi(x) = 2 \int_0^\infty \frac{dz z}{(z^2 + 1)^2} \left[ \frac{\sin(zx)}{zx} - 1 \right]$$

- It has an imaginary part ! (Laine, Philipsen, Romatschke, Tassler, 06; Escobedo, JS, 08; Brambilla, Ghiglieri, Vairo, Petreczky, 08)
- The dissociation temperature becomes

$$T_d \sim m_\mu \alpha^{2/3} / \ln^{1/3} \alpha < m_\mu \alpha^{1/2}$$

where  $m_\mu \alpha^{1/2}$  is the scale of the dissociation temperature for the screening mechanism (Matsui, Satz, 86)

# The $m\alpha/n \ll T \ll m$ case at $v \neq 0$

- $\Re V(r, T)$  calculated before (Chu, Matsui, 89)
- $V(r, T)$  is given by the Fourier transform of the longitudinal photon propagator  $\Delta_{11}(k)$  at  $k^0 = 0$  in the  $v$ -dependent HTL Lagrangian

$$\Delta_{11}(k) = \frac{1}{2}[\Delta_R(k) + \Delta_A(k) + \Delta_S(k)]$$

$\Delta_R^*(k) = \Delta_A(k)$ ,  $\Delta_S(k)$  contains the imaginary part

- $\Delta_S(k)$  must be calculated through the following formula, which differs from the one of the  $v = 0$  case (Carrington, Hou, Thoma, 97)

$$\Delta_S(k, u) = \frac{\Pi_S(k, u)}{2i\Im\Pi_R(k, u)} (\Delta_R(k, u) - \Delta_A(k, u))$$

$u = \gamma(1, \mathbf{v})$ . Recall that in the real time formalism  $\Pi_R = \Pi_{11} + \Pi_{12}$ ,  $\Pi_S = \Pi_{11} + \Pi_{22}$ ,  $\Delta_R(k) = 1/(\mathbf{k}^2 - \Pi_R)$

## The $m\alpha/n \ll T \ll m$ case at $v \neq 0$

- $\Pi_R(k, u)$  is a (complex) function of  $v$  and  $\theta$ ,  
 $\mathbf{k}\mathbf{v} = |\mathbf{k}|v \cos \theta$ , that reduces to  $-m_D^2$  when  $v = 0$
- We obtain

$$\Pi_S(k, u) = \frac{i2\pi m_D^2 T (1 - v^2)^{3/2} (1 + \frac{v^2}{2} \cos^2 \theta)}{|\mathbf{k}| (1 - v^2 \sin^2 \theta)^{5/2}}$$

# The $m\alpha/n \ll T \ll m$ case at $v \neq 0$

- The dissociation mechanism:

# The $m\alpha/n \ll T \ll m$ case at $v \neq 0$

• The dissociation mechanism:

- Let  $m_d$  be the momentum scale for which the real and imaginary part of the potential are equal. If  $m_d \gg m_D$ , as in the  $v = 0$  case, we find (for  $\theta \approx \pi/2$ )  $m_d \sim \alpha^{1/3} T \sqrt{1 - v^2}$ . Hence the dissociation temperature decreases. As long as  $m_D \ll m_d$ , the Landau damping remains the dominant mechanism

# The $m\alpha/n \ll T \ll m$ case at $v \neq 0$

## • The dissociation mechanism:

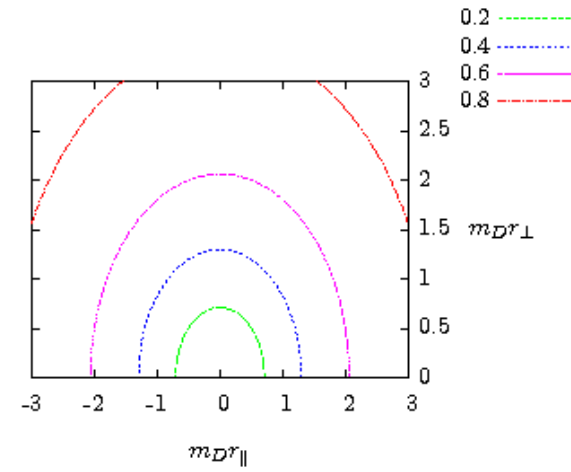
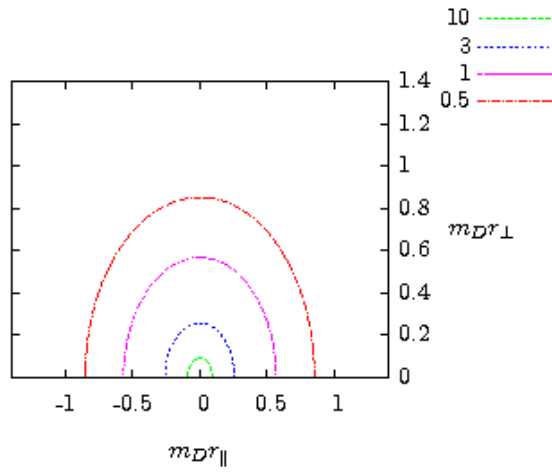
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- Since  $m_D$  remains non-vanishing when  $v \rightarrow 1$ , there will be some critical  $v$  for which  $m_d = m_D$  and both mechanisms Landau damping and screening compete

# The $m\alpha/n \ll T \ll m$ case at $v \neq 0$

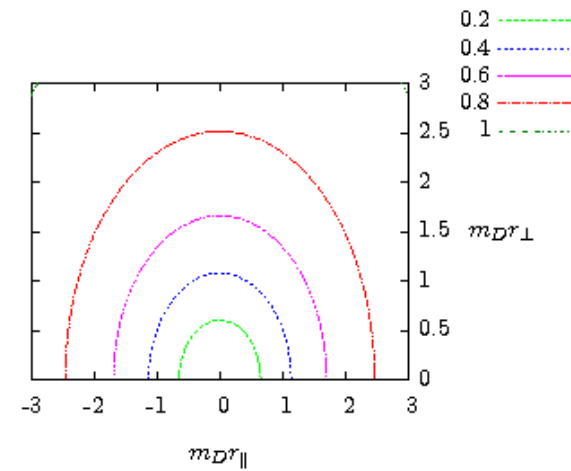
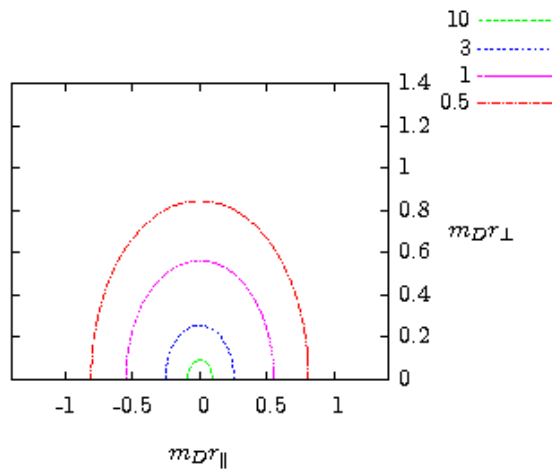
## • The dissociation mechanism:

- Let  $m_d$  be the momentum scale for which the real and imaginary part of the potential are equal. If  $m_d \gg m_D$ , as in the  $v = 0$  case, we find (for  $\theta \approx \pi/2$ )  $m_d \sim \alpha^{1/3} T \sqrt{1 - v^2}$ . Hence the dissociation temperature decreases. As long as  $m_D \ll m_d$ , the Landau damping remains the dominant mechanism
- Since  $m_D$  remains non-vanishing when  $v \rightarrow 1$ , there will be some critical  $v$  for which  $m_d = m_D$  and both mechanisms Landau damping and screening compete
- For  $v$  sufficiently close to 1 the imaginary part is negligible and the dominant mechanism for dissociation is screening

$$v = 0$$

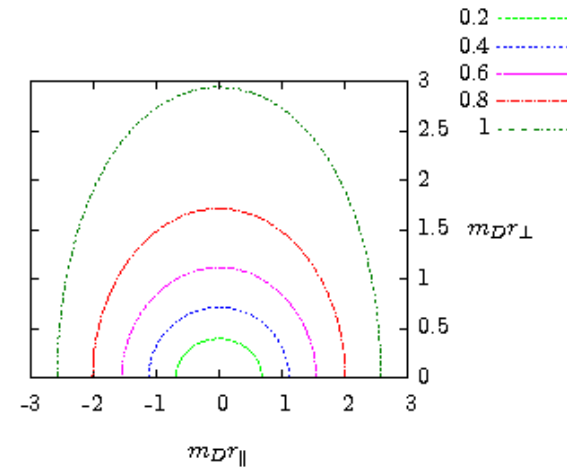
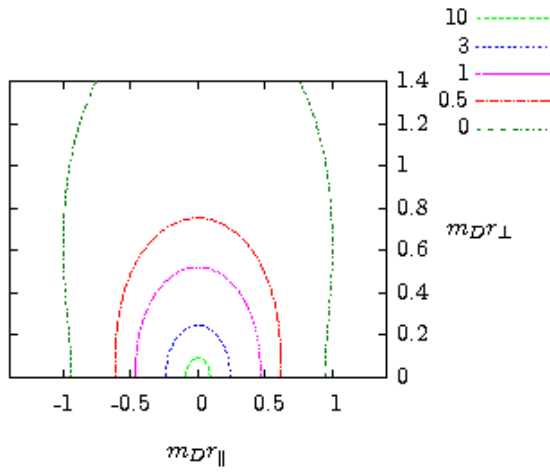


$$v = 0.5$$

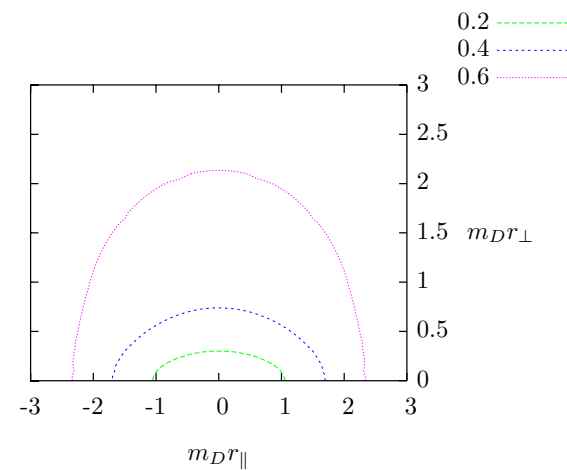
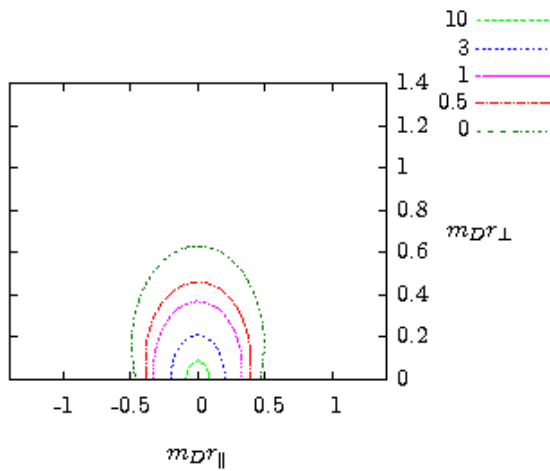




$$v = 0.9$$



$$v = 0.99$$





# Conclusions

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# Conclusions

- The decay width of a heavy quarkonium moving in a QGP decreases with its velocity

# Conclusions



- The decay width of a heavy quarkonium moving in a QGP decreases with its velocity
- The dominant mechanism for dissociation changes from Landau damping to screening as the velocity increases



