Open-charm Euclidean correlators within heavy-meson EFT interactions

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Received: 30 July 2020 / Accepted: 7 November 2020 / Published online: 20 November 2020
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Communicated by Ralf Rapp

Abstract The open-charm Euclidean correlators have been computed for the first time using the thermal spectral functions extracted from a finite-temperature self-consistent unitarized approach based on a chiral effective field theory that implements heavy-quark spin symmetry. The inclusion of the full-energy dependent open-charm spectral functions in the calculation of the Euclidean correlators leads to a similar behaviour as the one obtained in lattice QCD for temperatures well below the transition deconfinement temperature. The discrepancies at temperatures close or above the transition deconfinement temperature could indicate that higher-energy states, that are not present in the open-charm spectral functions, become relevant for a quantitative description of the lattice QCD correlators at those temperatures. In fact, we find that the inclusion of a continuum of scattering states improves the comparison at small Euclidean times, whereas differences still arise for large times.

1 Introduction

Understanding the medium modification of heavy mesons when embedded in a hot and/or dense matter is the subject of ongoing experiments, such as LHC and RHIC, whereas a great effort is also being devoted from the theoretical side (see [1–4] for reviews). Since heavy flavour (charm and beauty) is produced in the early stages of high-energy heavy-ion collisions (HiCs), heavy quarks and their hadronization into heavy mesons turn out to be excellent probes of the properties of the hot and dense medium created during the collisions. In fact, the suppression of quarkonium states, such as the \( J/\psi \) meson, in HiCs as compared to proton-proton collisions is widely considered as a signature of the deconfinement of hadronic matter into the quark-gluon plasma (QGP) [5]. This suppression could be also modified due to the in-medium changes of heavy mesons, as described by the comover scattering scenario (see, for example, Refs. [6–9]).

Lattice QCD is a powerful tool to study the in-medium modification of heavy mesons through the determination of their spectral properties in matter (see Ref. [10] for a recent review and references therein). Despite the recent progress in lattice QCD calculations, there are still a few drawbacks that prevent lattice results from being decisive when determining the spectral features of heavy mesons. From lattice QCD one can determine the so-called Euclidean meson correlators and the meson spectral functions are then extracted from them. The reconstruction of the spectral functions from the correlators turns out, however, to be rather complicated [11]. Furthermore, the simulation of light quarks on the lattice is computationally very demanding and usually larger (unphysical) masses are used.

Effective field theories in matter offer a complementary strategy to lattice QCD in order to determine the modification of the heavy-meson spectral features in a hot and/or dense medium. Matter below the deconfinement transition temperature consists of hadrons, essentially light mesons, in the low-density high-temperature regime. In this domain, the thermal properties of scalar and vector charm mesons have been recently obtained within a finite-temperature self-consistent unitarized approach based on a chiral effective field theory.
that implements heavy-quark spin symmetry [12, 13]. Once the spectral features are known, it is then possible to determine the corresponding Euclidean meson correlators and compare to lattice QCD results. In this way, the ill-posed extraction of the spectral function is avoided while testing directly the results from finite-temperature effective field theories against lattice QCD simulations.

In the present paper we determine the Euclidean meson correlators for open charm mesons and compare to the lattice QCD simulations of Ref. [14]. To the best of our knowledge, this work is the only computation of Euclidean correlators of open-charm mesons. We adapt our calculations of open-charm spectral functions in a pionic bath [12, 13] to the use of the unphysical masses determined in Ref. [14]. The paper is organized as follows. In Sect 2 we introduce the concept of the meson spectral function at finite temperature, while presenting in Sect 3 the calculation of the Euclidean correlators and spectral functions in lattice QCD. In Sect 4 we summarize our calculation for the open-charm spectral function within the effective field theory employing the meson unphysical masses reported in Ref. [14], whereas in Sect 5 we present our results for the Euclidean correlators and compare them with those from lattice QCD. Finally, in Sect 6 we give our conclusions and future outlook.

2 Spectral functions of mesons at finite temperature

The meson spectral function at finite temperature contains information not only on the mass and width of the ground state but also the masses and widths of the possible excited bound states as well as the continuum of scattering states. A schematic picture is shown in Fig. 1. At $T = 0$ the spectral function results from the contribution of different delta functions corresponding to the ground state of mass $m$ and the bound excited states, and a continuum distribution starting at $2m$ for 2-particle states. At finite temperature, one expects the masses to be modified as well as a broadening of 1-particle states to take place.

There exist a few theoretical approaches that can be used to determine the features of the meson spectral function, but none of them is yet conclusive in the full range of energies and temperatures available in the experiments. Perturbative QCD can be only applied at very large energies and/or temperatures [15–17]. Using lattice QCD, the meson correlator can be calculated from first principles for any energy and temperature (a priori), but the spectral function needs to be reconstructed from the lattice data, which is a non-trivial task [10, 11]. The AdS/CFT duality has also been used to describe some features of the spectral function, but a clear correspondence with QCD allowing quantitative studies is still missing [18, 19]. Hadronic models based on effective theories at finite temperature are an additional tool to learn about the spectral function below the deconfining temperature [1–4]. Therefore, only the interplay between these techniques may shed light on this issue.

3 Euclidean correlators and spectral functions in lattice QCD

The primary tools in lattice QCD calculations are Euclidean correlators of some operators $\hat{O}$, described by the path integral over all degrees of freedom:

$$\langle \hat{O}_1(\tau, \mathbf{x}) \hat{O}_2(0, 0) \rangle = \frac{1}{Z} \int \mathcal{D}[\bar{\psi}, \psi] \times \hat{O}_2[U, \bar{\psi}, \psi] \hat{O}_1[U, \bar{\psi}, \psi] e^{-S_F[U, \bar{\psi}, \psi] - S_G[U]},$$

(1)

with $Z$ being the partition function

$$Z = \int \mathcal{D}[\bar{\psi}, \psi] e^{-S_F[U, \bar{\psi}, \psi] - S_G[U]},$$

(2)

where $S_F[U, \bar{\psi}, \psi]$ is the fermion part and $S_G[U]$ the gluon part of the discretized QCD action.

In the case of a meson (quark anti-quark pair) the operator to consider is $J_H(\tau, \mathbf{x}) = \bar{\psi}(\tau, \mathbf{x}) \gamma_\mu \psi(\tau, \mathbf{x})$, where $\gamma_\mu = 1, \gamma_5, \gamma_\mu, \gamma_5 \gamma_\mu$ correspond to the scalar, pseudoscalar, vector and axial vector channels, respectively, and $f$ refers to the flavour of the valence quark.

In the case of meson correlators in a hot medium one has to consider that the temperature of the system on a lattice of size $N_a^3 \times N_t$, with lattice spacing $a$, is related to the temporal extent through $T = 1/(\alpha N_t)$. The Euclidean temporal correlator in momentum space

$$G_E(\tau, \mathbf{p}; T) \equiv \int d^3x e^{-ip \cdot x} \langle J(\tau, \mathbf{x}) J(0, 0) \rangle$$

(3)
is related to the spectral function $\rho(\omega, p; T)$ through the convolution with a known kernel $K(\tau, \omega; T)$:

$$G_E(\tau, p; T) = \int_0^\infty d\omega K(\tau, \omega; T)\rho(\omega, p; T),$$

(4)

with

$$K(\tau, \omega; T) = \frac{\cosh[\omega(\tau - 1/T)]}{\sinh(\omega/T)}.$$  

(5)

In the following we omit the dependence on the momentum $p$, as we focus on spectral functions with $p = 0$ for simplicity, and make the identifications $G_E(\tau; T) \equiv G_E(\tau; p = 0; T)$ and $\rho(\omega; T) \equiv \rho(\omega, p = 0; T)$.

In lattice QCD simulations values of the Euclidean correlator are obtained for a set of points in Euclidean time, $\tau = \tau_i$, i.e. \{ $\tau_i, G_E(\tau_i; T)$ \} for $i = 1, ..., N_\tau$ and $\tau_i \in [0, 1/T]$. In addition, the lattice data $G_E(\tau_i; T)$ have a statistical error due to the fact that only a finite number of gauge configurations can be generated in a Monte Carlo simulation. The inversion of Eq. (4) to extract a continuous spectral function $\rho(\omega; T)$ from such limited number of data points with noise is an ill-posed problem. Two methods are usually employed to try to circumvent this problem, both taking specific assumptions on the shape of the spectral function: (i) Bayesian methods like the maximum entropy method (MEM) or stochastic reconstruction methods [20], which perform the kernel inversion by statistically inferring the most probable spectral function; and (ii) fitting the lattice data with suitable Ansätze for the spectral function, incorporating bound states, a continuum and perturbative input [21]. The fact that a priori assumptions about the spectral functions are needed makes the determination of their shape and details at finite temperature very challenging, as we do not have much prior information on them.

By inspecting Eq. (4), one can see that the temperature dependence of the correlators not only comes from the spectral function but also from the inherent temperature dependence of the integration kernel. When directly comparing correlation functions at different temperatures, one may want to discern the differences due to the modification of the spectral function with temperature alone. In order to do so, it is useful to define the so-called reconstructed correlator at a reference temperature $T_r$,

$$G'_E(\tau; T, T_r) = \int_0^\infty d\omega K(\tau, \omega; T)\rho(\omega; T_r).$$

(6)

The integration kernel is the same as that of $G_E(\tau; T)$ and therefore any difference when comparing $G_E(\tau; T)$ and $G'_E(\tau; T, T_r)$ arises from differences in the spectral functions at $T$ and $T_r$. The value of $T_r$ is usually chosen to correspond to a temperature at which the shape of the spectral function is better known, so one can reliably trust the spectral function obtained from the lattice correlator. Thus, the lowest temperature available is usually chosen.

In this paper we use the lattice setup of Ref. [14], where an anisotropic lattice with spacing in the temporal direction $a^{-1}_t = 5.63$ GeV and anisotropy parameter $\xi = 3.5$ is used. The ensembles employed contain dynamical light and strange quarks, with unphysical masses for the two mass-degenerate light quarks and roughly physical values for the strange and charm quarks. The resulting masses of the light and charm mesons on the lattice are the following: $m_\pi = 384$ MeV, $m_K = 546$ MeV, $m_\eta = 589$ MeV, $m_D = 1880$ MeV, $m_{Ds} = 1943$ MeV. The pseudocritical temperature determined for this configuration is $T_c = 185$ MeV and the correlators have been calculated for temperatures $T = (44, 141, 156, 176, 201, 235, 281, 352)$ MeV. For further details, we refer the reader to Ref. [14] and references therein.

### 4 Spectral functions within an effective theory at finite temperature

In Refs. [12,13] we have obtained the spectral functions of the $D$ and $D_s$ mesons in a pionic bath at finite temperature. We use an effective Lagrangian describing the interactions of open heavy-flavour mesons ($D$ and $D_s$) with the light pseudoscalar mesons ($\pi, K, \bar{K}$ and $\eta$) that is based on chiral and heavy-quark spin-flavour symmetries. The details of the interaction Lagrangian, as well as the different coupled channels in each of the strangeness ($S$) sectors of the resulting charm meson - light meson interactions, can be found in Ref. [13]. The parameters of the model are taken from the study of Ref. [22], where they are fitted to reproduce the charm-meson masses and the scattering lengths from lattice simulations. In the present work we employ the amplitudes of the channels with $S = 0$ and 1, as they are the ones needed, respectively, to describe the properties of the $D$ and $D_s$ mesons in a pionic medium. These amplitudes have been obtained from a unitarization procedure based in solving the Bethe-Salpeter integral matrix equation in a full coupled-channels basis

$$T_{ij} = V_{ij} + V_{ik} G_{D\Phi,k} T_{kj}$$

(7)

consisting in a resummation of an infinite number of multiple scattering diagrams. In Eq. (7), $T$ is the scattering amplitude, $V$ is the interaction potential extracted from the effective Lagrangian and $G_{D\Phi}$ is the meson-meson loop function. The subindices $i, j, k$ denote the incoming and outcoming charm meson($D$) - light meson($\Phi$) channels. After a commonly employed factorization procedure, the integral Eq. (7) becomes purely algebraic and can be easily inverted to obtain
the scattering amplitude $T_{ij}$. The loop function $G_{D\Phi}$ is divergent and needs to be properly regularized by means of a cut-off or using dimensional regularization. As discussed in detail in Ref. [13], we employ the former method, with a naturally sized cut-off of $\Lambda = 800$ MeV, because it can be extended to finite temperature in a straightforward way.

The effects of a medium at finite temperature are introduced in this model within the imaginary time formalism (ITF), in which the energy integration is replaced by a sum over the so-called Matsubara frequencies, and the mesons in the loop function are dressed with the spectral functions. After the Matsubara summation and the analytical continuation to real energies, the loop function entering the Bethe-Salpeter equation reads:

$$G_{D\Phi}(E, \mathbf{p}; T) = \int \frac{d^3q}{(2\pi)^3} \int d\omega \int d\omega' \frac{S_D(\omega, \mathbf{q}; T)S_D(\omega', \mathbf{p} - \mathbf{q}; T)}{E - \omega - \omega' + i\epsilon}[1 + f(\omega, T) + f(\omega', T)],$$

where $f(\omega, T)$ is the Bose-Einstein distribution for bosons. We note that here and in Refs. [12,13] we have neglected the thermal modification of the pions since their change with temperature is expected to be small [13]. Thus vacuum spectral functions are taken for the light mesons.

The spectral function of the charm meson is defined in terms of its propagator in a hot medium, $D_D$, as

$$S_D(\omega, \mathbf{q}; T) = -\frac{1}{\pi} \text{Im} D_D(\omega, \mathbf{q}; T) = -\frac{1}{\pi} \text{Im} \left( \frac{1}{\omega^2 - \mathbf{q}^2 - M_D^2 - \Pi_D(\omega, \mathbf{q}; T)} \right),$$

where the self-energy, $\Pi_D$, is obtained from closing the pion line in the $T_{D\pi \rightarrow D\pi}$ matrix element of the unitarized amplitude. Heavier light meson contributions (e.g. kaons) to the dressing of the $D$ meson are neglected as their abundance in the hot medium is suppressed. Indeed in Ref. [13] we have quantified the effect of the kaons in the medium to be around 10% of the mass shift of the $D$ mesons induced by the pion-only hot medium at $T = 150$ MeV. In the ITF the pion contribution to the $D$-meson self-energy is given by

$$\Pi_D(\omega, \mathbf{q}; T) = \int \frac{d^3q'}{(2\pi)^3} \int dE \frac{f(E, T) - f(\omega_n, T)}{\omega_n \omega^2 - (\omega_n - E)^2 + i\epsilon} \times \left( \frac{1}{\pi} \right) \text{Im} T_{D\pi \rightarrow D\pi}(E, \mathbf{p}; T),$$

where $\mathbf{q}' = \mathbf{p} - \mathbf{q}$.

The set of Eqs. (7)–(10) needs to be solved iteratively until convergence is achieved to ensure the self-consistency of the results. In this hadronic effective model the masses of the mesons in vacuum are input parameters and, therefore, can be tuned to those used in lattice QCD for a straightforward comparison of both approaches.

The spectral functions of the $D$ and $D_s$ mesons obtained from their interaction with the unphysical heavy pions in a hot medium are shown in Fig. 2. In these figures we show the spectral functions for $D$ (upper panel) and $D_s$ (lower panel) as a function of the meson energy for the different temperatures used in Ref. [14]. As discussed in Refs. [12,13], we clearly see the increased broadening of both spectral functions with temperature due to the larger available phase space for decay at finite temperature. Moreover, the maximum of both spectral functions slightly moves to lower energies with temperature given the attractive character of the heavy meson - light meson interaction.

A study of net charm fluctuations [23] suggests that open charm hadrons start to dissolve already close to the chiral crossover. Although the deconfined degrees of freedom necessary to investigate the melting of charm hadrons are absent in the model, we still show the spectral functions coming

![Fig. 2](image-url) Spectral functions of the $D$-meson (upper panel) and $D_s$-meson (lower panel) obtained from their effective interaction with unphysically heavy pions at finite temperature.

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from our hadronic model for temperatures above and close to the lattice pseudocritical temperature $T_c$ in order to explore the validity of our results for the correlators at those temperatures, as data below $T_c$ is scarce.

5 Results of Euclidean correlators and comparison with lattice QCD

Once the $D$ and $D_s$ spectral functions at finite temperature are known, we can obtain the corresponding Euclidean correlators from Eqs. (4) and (6) so as to compare them to lattice QCD calculations. It is important to notice that the spectral function $S_D(\omega; T)$ of Eq. (9), displayed in Fig. 2, differs from $\rho(\omega; T)$, entering in Eq. (4). This is due to the fact that the former contains the ground-state peak and an additional continuum, corresponding to $D\pi\pi$ scattering states in the case of the $D$ spectral function and $D\pi K$ states in the case of the spectral function of the $D_s$, while the latter contains all possible quark-antiquark ($c\bar{c}$ or $c\bar{s}$) states.

Furthermore, the dimensions of $S_D(\omega; T)$ are MeV$^{-2}$ while the dimensions of $\rho(\omega; T)$ are MeV$^2$. They are related through the fourth power of the charm meson mass [24,25]:

$$
\rho_{gs}(\omega; T) = M_D^4 S_D(\omega; T),
$$

(11)

with $M_D$ the vacuum value in the PDG for the mass of the $D$ (or $D_s$) meson. In Refs. [24,25] the authors obtained the relation between the electromagnetic current-current correlation in matter and the vector meson self-energy, and hence the meson spectral function, based on vector-meson dominance. In Eq. (11) we have similarly connected the Euclidean current-current correlator given by the lattice simulation with the open-charm spectral function resulting from the chiral effective theory that implements heavy-quark spin symmetry. Note also that in lattice QCD studies the reconstructed spectral function is usually identified with the dimensionless quantity $\rho(\omega; T)/\omega^2$.

With $\rho_{gs}(\omega; T)$ and $\rho_{gs}(\omega; T)$ we can readily calculate Euclidean correlators and the ratios with the reconstructed correlators, $G_E(\tau; T)/G_E(\tau; T)$, for a direct comparison with lattice data. The input spectral function at $T = (141, 156, 176, 201, 235)$ MeV and $T_r = 44$ MeV are shown with solid lines in the left panel of Fig. 3 for the $D$ meson and in the left panel of Fig. 5 for the $D_s$ meson, while the Euclidean correlators are displayed in solid lines in the corresponding right panels together with the lattice data of Ref. [14] (coloured filled circles). Also, in the left panels of Figs. 4 and 6 the ratios for the correlators are shown with solid lines together with the lattice results of Ref. [14] (coloured filled circles with error bars).

The first clear observation is the deviation of the correlators (solid lines in the right panels of Figs. 3 and 5) and the ratio of correlators (left panels of Figs. 4 and 6) at small Euclidean times $\tau$ with respect to the lattice data points for all the temperatures and for both the $D$ and $D_s$ mesons. However, we note that for the lowest temperature $T = 141$ MeV, the ratio of correlators lies within the error bars of the lattice data. For increasing temperatures, the calculated ratios deviate largely from the lattice calculations. Above or close to the pseudocritical temperature, $T_c = 185$ MeV, we do not expect a good matching as the deconfined degrees of freedom are not included in the hadronic model described above, but one would expect a better comparison at lower temperatures.

The discrepancy observed at small $\tau$ for temperatures below $T_c$ might be due to the fact that the spectral functions do not contain the higher-energy states present in the lattice correlators in addition to the ground-state, i.e. possible excited states and the continuum spectrum. As a first approximation, in the following we only add a continuum contribution to the spectral functions. In this way, we aim at understanding the differences with the fewest possible parameters, while trying to improve the comparison of the hadronic and lattice approaches.

With this goal we define the lattice spectral function as

$$
\rho(\omega; T) = \rho_{gs}(\omega; T) + a \rho_{cont}(\omega; T),
$$

(12)

where we add, to the ground-state spectral function obtained from the effective field theory, the contribution of a continuum of scattering states weighted with a factor $a$.

The continuum contribution to the spectral function is sometimes mimicked with a step function. A parametrization of the free meson spectral function in the non-interacting limit is also often used and was first derived for charmonium states in Refs. [15,26]. This spectral function describes quark-antiquark pairs with degenerate masses in the limit of infinitely high temperature. An equivalent expression can be derived in the case of non-degenerate quark masses $m_1 > m_2$ [27].

$$
\rho_M(\omega; T) = \frac{N_c}{32\pi} \left( \frac{m_1^2 - m_2^2}{\omega^2} + 1 \right) \frac{2m_1^2}{\omega^2} \omega^2
\times \left[ (a_M - b_M) + \frac{2m_1^2 + m_2^2}{\omega^2} - 4c_M \frac{m_1 m_2}{\omega^4} \right]
\times [n(-\omega_0, T) - n(\omega - \omega_0, T)] \theta (\omega - (m_1 + m_2)),
$$

(13)

where $N_c = 3$ is the numbers of colours, $\omega_0 = \frac{1}{2\omega}(\omega^2 + m_1^2 - m_2^2)$, $n(\omega, T) = [e^{\omega/T} + 1]^{-1}$ is the Fermi-Dirac distribution and the coefficients $(a_M, b_M, c_M)$ are $(1, -1, 1)$ for the scalar, $(1, -1, -1)$ for the pseudoscalar, $(2, -2, -4)$ for the vector, and $(2, -2, 4)$ for the pseudovector channels.
Therefore, for pseudoscalar mesons we have

\[
\rho_{\text{cont}}(\omega; T) = \frac{3}{32\pi} \left( \frac{m_1^2 - m_2^2}{\omega^2} + 1 \right)^2 \frac{4m_2^2}{\omega^2} \omega^3 
\times 2 \left( 1 - \frac{(m_1 - m_2)^2}{\omega^2} \right) \times [n(-\omega_0, T) - n(\omega - \omega_0, T)] \theta(\omega - (m_1 + m_2)).
\]

In the case of the $D$ meson we take $m_1 = m_c = 1.5$ GeV and $m_2 = m_l = 0$, whereas for $D_s$ we use $m_1 = m_c = 1.5$ GeV and $m_2 = m_s = 100$ MeV. The left panels of Figs. 3 and 5 contain the spectral functions obtained for three different values of the continuum to ground-state contribution: $a = 0$ (solid lines, no continuum), $a = 1$ (dashed lines) and $a = 10$ (dotted lines), for the $D$ and $D_s$, respectively. The corresponding Euclidean correlators are plotted in the right panel.
panels of Figs. 3 and 5 and the ratios with the reconstructed correlators are shown in Figs. 4 and 6.

The inclusion of the continuum in the spectral functions improves the behaviour of the correlators and the ratio of correlators at small $\tau$, but does not allow to reproduce the shape of the lattice correlators\(^\dagger\). It also permits the ratios to go to one at $\tau \to 0$ for all temperatures, as the region of very small Euclidean times is essentially governed by the spectral function at very high energies. However, the modification of the ratios at larger $\tau$ due to the inclusion of the continuum is rather moderate and only the results for the lowest temperature of $T_c = 141$ MeV are compatible with the lattice data within the error bars. The region of middle and large values of $\tau$ is rather sensitive to the shape of the spectral functions at low energies (few GeV), where not only varia-

\(^\dagger\) Note that the lattice correlators we used here are not continuum extrapolated and therefore suffer from cut-off effects at small $\tau \lesssim 0.1$ fm.
tions in the ground state properties can produce significant modifications of the correlators but also where the free spectral function might be not enough to describe the continuum and where excited states are likely to be present.

In particular the behaviour of the correlators can vary due to a widening of the ground state, which is expected at temperatures close to $T_c$ when considering further contributions to the $D$ and $D_s$ meson self-energies coming from the thermal kaons and antikaons in the medium. In our previous work [13] we have shown that the inclusion of $K$ and $\bar{K}$ mesons in the bath in addition to pions starts to have an appreciable effect to the mass shift of the $D$ and $D_s$ mesons at temperatures $T \sim 140$–150 MeV. Yet the impact on the $D_s$ width is visible already at $T = 80$ MeV and the $D_s$ width in a kaonic and pionic medium is almost a factor three larger than that induced in a pion-only medium at $T = 150$ MeV. In the present paper unphysical masses for pions and kaons are used, with the pions being almost three times heavier than physical pions and the kaons similarly heavy, thus considerably reducing the mass gap between the two lightest pseudoscalar mesons. As a consequence, the temperature onset for a sizable thermal modification of the heavy mesons due to pions is larger and closer to that of kaons, in such a way that the kaon-induced width of the charm mesons, specially of the $D_s$, might not be small at temperatures close to $T_c$. Therefore, considering a pionic and kaonic medium in our calculations could improve the comparison of the ratios of Euclidean correlators of the $D_s$ meson at $T = 141$ MeV and of both $D$ and $D_s$ mesons at larger temperatures below $T_c$.

In what regards the excited states, their properties and thermal modification are not known and their inclusion in the spectral function is not feasible. An alternative way to eliminate this source of discrepancy when comparing the results of the calculated Euclidian correlators with those simulated on a lattice could be to apply techniques for the reduction of excited-state contamination from the correlators in the analysis of the lattice QCD data.

Finally, one might also wonder about the validity of the chiral effective theory employed at an unphysical pion mass of 384 MeV. However, we consider the Lagrangian to be reliable, as its parameters have been adjusted to finite volume energy levels and scattering lengths, obtained at various unphysical masses but having a smooth extrapolation down to the physical pion mass [22].

6 Conclusions and outlook

In this work we have computed for the first time Euclidean correlators for the charm $D$ and $D_s$ mesons from their corresponding thermal spectral functions obtained within a finite-temperature self-consistent unitarized approach based on a chiral effective field theory that implements heavy-quark spin symmetry. The goal is to compare the calculated correlators with those obtained in lattice QCD simulations at unphysical masses.

We have started by analyzing the behaviour of the Euclidean correlators once the full energy-dependent spectral functions are considered and have found that for temperatures well below the deconfinement transition temperature the ratio of correlators lies within the error bars of the lattice data. However, as we increase the temperature, the ratios deviate significantly from the lattice predictions due to the fact that the spectral functions do not contain the higher-energy components of the bound excited states as well as the continuum of scattering states that are present in the lattice correlators in addition to the ground-state.

We have therefore studied the addition of a continuum contribution to the spectral functions, so as to determine its effect on the Euclidean correlators at finite temperature, while aiming at improving the comparison between the hadronic and lattice QCD calculations. The inclusion of the continuum in the spectral functions improves the behaviour of the correlators at small $\tau$ for all the temperature studied, but it is not sufficient to reproduce the shape of the lattice correlators for temperatures close or above the deconfinement transition temperature. Indeed, the region of intermediate to large $\tau$ is very sensitive to the possible existence of other excited states ignored in our model.

In the future we plan to improve the comparison with lattice QCD results for the correlators by considering finite cutoffs in our calculations. Moreover, we will also study the effect of including the thermal modification of the charm ground-state meson properties induced by a hot kaonic medium. We also aim at computing the Euclidean correlators using physical masses, so as to predict the expected behaviour of the lattice QCD correlators for temperatures below the deconfinement transition, where our effective approach is fully justified.

Acknowledgements The authors thank J.I. Skullerud and A. Rothkopf for kindly providing the lattice QCD data. G.M. and A.R. acknowledge support from the Spanish Ministerio de Economía y Competitividad (MINECO) under the project MDM-2014-0369 of ICCUB (Unidad de Excelencia “María de Maeztu”), and, with additional European FEDER funds, under the contract FIS2017-87534-P. G.M. also acknowledges support from the FPU17/04910 Doctoral Grant from the Spanish Ministerio de Educación, Cultura y Deporte and from a STSM Grant from THOR COST Action CA15213. L.T. acknowledges support from the FPA2016-81114-P Grant from the former Ministerio de Ciencia, Innovación y Universidades, the PID2019-110165GB-I00 Grant from the Ministerio de Ciencia e Innovación, the Heisenberg Programme of the Deutsche Forschungsgemeinschaft (DFG, German research Foundation) under the Project Nr. 383452351 and Nr. 411563442 (Hot Heavy Mesons), and the THOR COST Action CA15213. L.T. and O.K. acknowledge support from the DFG through the Grant Nr. 315477589 - TRR 211 (Strong-interaction matter under extreme conditions). We also thank the EU STRONG-2020 project under the program H2020-INFRAIA-2018-1, grant agreement no. 824093.
Data Availability Statement This manuscript has no associated data or the data will not be deposited. [Authors’ comment: The data used in this paper are available upon request via email.]

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