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$T_{cc}(3875)$ at finite density and finite temperature

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We investigate how the $T_{cc}(3875)$ and its charge-conjugated partner behave when they are embedded in a dense nuclear medium or in a hot pionic medium. In vacuum, this state is modeled as an isoscalar DD^* S-wave bound state, formed through effective interactions that correspond to different Weinberg compositeness models. Medium effects are introduced via the two-meson loop function, incorporating the self-energies acquired by $D^{(*)}$ mesons when they interact with the particles in the dense or hot mediums. The obtained lineshapes in the finite density nuclear matter reveal a charge-conjugation asymmetry stemming from the very different $D^{(*)}N$ and $\overline{D}^{(*)}N$ interactions. The temperature-induced changes on the T_{cc} do not exhibit this charge-conjugation asymmetry, and are predicted to be extremely large near the crossover temperature of $T = 150$ MeV. These predictions for both the finite-density and finite-temperature scenarios are highly sensitive to the Weinberg molecular component in the T_{cc} wave function. If these finite-density or finite-temperature patterns could be measured, they |would provide additional insight into the inner structure of this exotic state.

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PoS(HADRON2025)133

1. Introduction

The doubly-charmed $T_{cc}(3875)$ state discovered by LHCb in 2021 [1] is one of the clearer signals of a hadron resonance with an exotic structure, since it needs a minimum quark content of four quarks. This very narrow state is observed in the $D^0D^0\pi^+$ invariant–mass spectrum. Its mass lies extremely close to the $D^{*+}D^0$ threshold, with a binding energy $\delta m_{\rm exp} = m_{T_{cc}} - m_{\rm thr}$ of $-360 \pm 40^{+4}_{-0}$ keV, where $m_{\rm thr} = 3875.09$ MeV is the $D^{*+}D^0$ threshold. The measured width is $\Gamma = 48 \pm 2^{+0}_{-14}$ keV [2]. The T_{cc} has been and continues to be of great interest in the hadronic community. There exist a great number of theoretical studies based on different assumptions about the internal structure of this tetraquark. Some of these descriptions are based on hadron-molecular models — this interpretation being supported by the proximity of this state to the $D^{*+}D^0$ and $D^{*0}D^+$ thresholds — while others lean towards the compact-tetraquark interpretation. In any case, the proximity of this state to the aforementioned two-meson thresholds makes it necessary to consider these hadronic degrees of freedom in order to correctly analyze the experimental data [3].

In this work, we employ a heavy-meson EFT to study the behavior of the T_{cc} when produced in dense nuclear matter or at high temperature. The low-energy constants of the effective D^*D interaction are tuned to reproduce different Weinberg–compositeness scenarios [4, 5], enabling a discussion in terms of the D^*D molecular component of the T_{cc} wave function. If the T_{cc} were to be mainly of molecular nature, we anticipate pronounced in-medium modifications of the T_{cc} width — potentially leading to the state's dissolution — due to the its small binding energy. If, instead, the T_{cc} is predominantly compact, the effects of the nuclear medium or temperature may be smaller and, in any case, qualitatively different from those discussed here.

2. Formalism

In this section, we briefly review the formalisms used in Refs. [6, 7] for the finite nuclear density and finite temperature scenarios. Working in the isospin limit, we model $T_{cc}(3875)^+$ as an S-wave DD^* state with quantum numbers $I(J^P) = 0(1^+)$. We then consider two families of energy-dependent contact interactions, expanded around threshold as:

$$V_A(s) = C_1 + C_2\,[s - (m_D + m_{D^*})^2]\,, \tag{1a}$$

$$V_B(s) = \left(\tilde{C}_1 + \tilde{C}_2[s - (m_D + m_{D^*})^2]\right)^{-1}\,. \tag{1b}$$

In this expression, the quantity s is the invariant mass squared of the two-meson system. m_D and m_{D^*} are the masses of the D and D^* mesons, respectively, taken as the average of the two components of the corresponding isodoublet. C_1 and C_2 ($\tilde{C}_1$ and $\tilde{C}_2$) are two low-energy constants.

In vacuum, the $T-$matrix is unitarized using the on-shell Bethe-Salpeter equation (BSE)

$$\mathcal{T}^{-1}_{A(B)}(s) = V^{-1}_{A(B)}(s) - \Sigma_0(s), \tag{2}$$

where Σ_0 is the two-meson loop function in vacuum, given by

$$\Sigma_0(s) = i\int \frac{d^4q}{(2\pi)^4}\frac{1}{(P-q)^2 - m_D^2 + i\varepsilon}\frac{1}{q^2 - m_{D^*}^2 + i\varepsilon}. \tag{3}$$

This divergent quantity needs to be regularized, for which we employ a cutoff $\Lambda = 700$ MeV in the three-momentum integration.

PoS(HADRON2025)133

In our model, the T_{cc} in vacuum appears as a DD^* bound state, with a binding energy of -800 keV.[1] Then, we impose the potential to be such that the unitarized $T-$matrix presents a pole at $s = m_0^2$, with m_0 the vacuum mass of the T_{cc}. We also fix the coupling g_0^2 of the T_{cc} to the DD^* channel, which is closely related to the compositeness P_0 [5]. These two conditions read

$$\mathcal{T}^{-1}(m_0^2) = 0, \tag{4a}$$

$$\left[\mathcal{T}^{-1}\right]'(m_0^2) = g_0^{-2} = P_0\Sigma_0'(m_0^2), \tag{4b}$$

where the prime ($\prime$) symbol denotes a derivative with respect to s; and they determine the values of the low-energy constants C_1 and C_2 (or $\tilde{C}_1$ and $\tilde{C}_2$).

When considering medium modifications, we solve the BSE using a modified two-meson loop function, leaving the interaction kernel $V_{A(B)}$ unaffected. In the *finite nuclear-density case*, the $D^{(*)}$ mesons are dressed with the nucleons of the Fermi sea using an $SU(8)$ spin-flavor interaction, as described in Ref. [8]. Once the spectral functions of the $D^{(*)}$ mesons have been determined, one can compute the modified loop function as

$$\Sigma(s;\rho) = \int_\Lambda \frac{d^3q}{(2\pi)^3} \int_0^\infty d\omega \int_0^\infty d\omega' \times \left(\frac{S_D(\omega,\vec{q};\rho) S_{D^*}(\omega',\vec{q};\rho)}{E-\omega-\omega'+i\varepsilon} - \frac{S_{\overline{D}}(\omega,\vec{q};\rho) S_{\overline{D}^*}(\omega',\vec{q};\rho)}{E+\omega+\omega'-i\varepsilon} \right). \tag{5}$$

Actually, due to the very different $D^{(*)}N$ and $\overline{D}^{(*)}N$ interactions (here N represents a nucleon), the charge conjugated partner of the $T_{cc}(3875)^+$, namely the $T_{\bar{c}\bar{c}}(3875)^-$, will get different density-dependent modifications. In vacuum, the description of the $T_{\bar{c}\bar{c}}^-$ is equivalent to that of the T_{cc}^+. However, when considering the nuclear medium, the $\overline{D}\,\overline{D}^*$ loop is modified in a different way, given by exchanging the particle and antiparticle spectral functions in the expression of Eq. (5).

Considering now the *finite temperature* scenario, the modifications to the T_{cc}, which again enter through the two-meson loop function, come from two sources: 1) the statistical distribution of the different hadron species (Bose-Einstein for mesons) and 2) the interactions of the $D^{(*)}$ mesons with the different species that populate the thermal bath (encoded in the temperature-dependent spectral functions). The modified loop function now reads as follows.

$$\begin{aligned}\Sigma(s;T) = \int \frac{d^3q}{(2\pi)^3} \int_0^\infty d\omega \int_0^\infty d\omega' S_D(\omega,\vec{q}\,;T) S_{D^*}(\omega',\vec{q}\,;T) \\ \times \Bigg\{ [1+f(\omega,T)+f(\omega',T)] \left(\frac{1}{E-\omega-\omega'+i\varepsilon} - \frac{1}{E+\omega+\omega'+i\varepsilon} \right) \\ + [f(\omega,T)-f(\omega',T)] \left(\frac{1}{E+\omega-\omega'+i\varepsilon} - \frac{1}{E-\omega+\omega'+i\varepsilon} \right) \Bigg\}.\end{aligned} \tag{6}$$

In this expression, the f functions stand for Bose-Einstein distributions, while S are the temperature-dependent spectral functions, obtained in Refs. [9, 10] from the $D^{(*)}$ interactions with a thermal bath of pions. In contrast to the nuclear-medium case, the temperature-dependent spectral functions of particle and antiparticle are the same.

[1] We do not take the physical binding energy since we are working in the isospin limit. In addition, we disregard the width of the T_{cc} in the vacuum, which is negligible compared to the in-medium width.

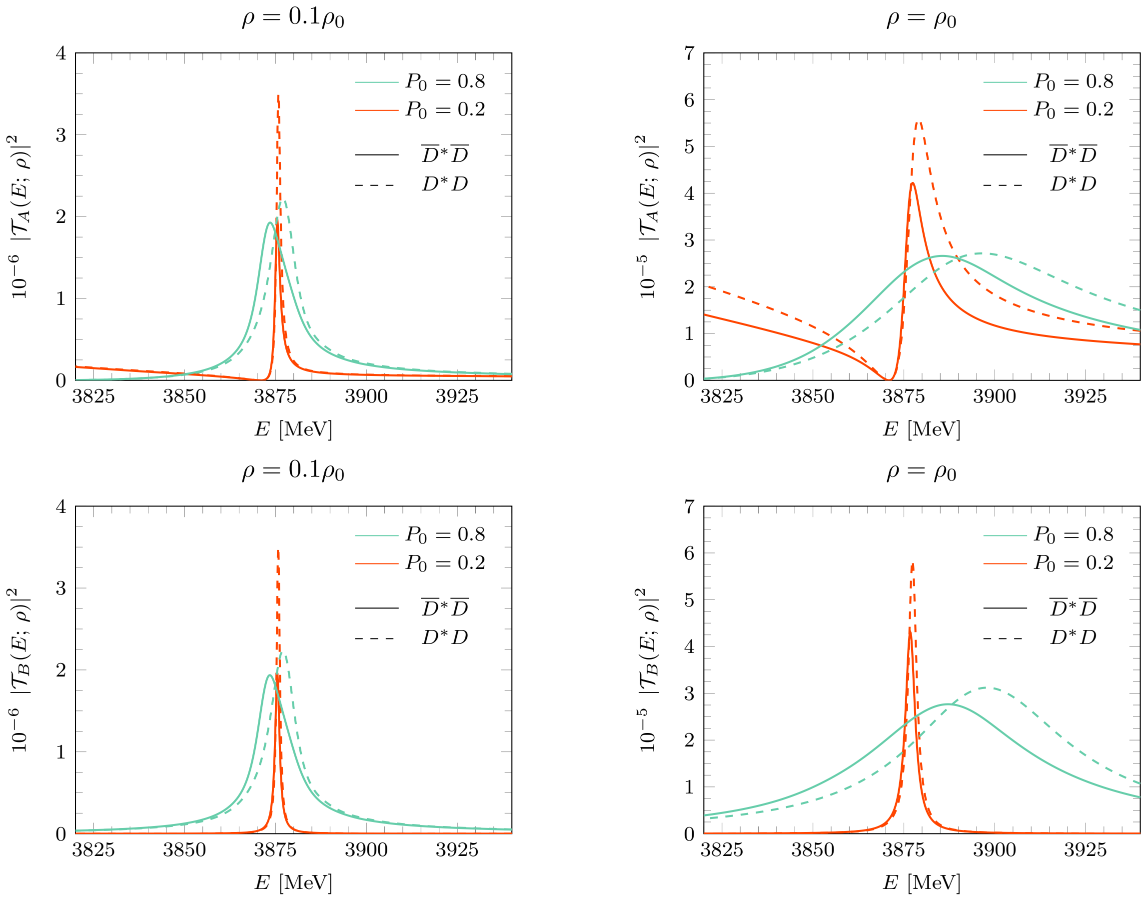

Figure 1: $\overline{D}^*\overline{D}$ (solid lines) and D^*D (dashed lines) modulus square amplitudes in nuclear matter obtained by solving the on-shell BSE using the $V_A(s)$ (top panels) and $V_B(s)$ (bottom panels) potentials, for vacuum molecular probabilities $P_0 = 0.2$ (dark orange) and $P_0 = 0.8$ (ligh green), and for different nuclear densities ρ in units of the normal nuclear density $\rho_0 = 0.17$ fm^{-3} (left and right columns).

3. Results and discussion

In this section, we show results for the DD^* (and $\overline{D}\overline{D}^*$) $I(J^P) = 0(1^+)$ amplitude — where the T_{cc} pole in vacuum shows up — in both the nuclear medium and at finite temperature. Let us start by discussing the nuclear density case, shown in Fig. 1. In all panels of this figure, the BSE is solved with the appropriate density–dependent two–meson loop, using kernel type A (top panels) or type B (bottom panels). We consider two densities — $0.1\rho_0$ (left column) and ρ_0 (right column), with $\rho_0 = 0.17$ fm^{-3} the normal nuclear density — and for each density two molecular probabilities, $P_0 = 0.2$ (dark orange) and $P_0 = 0.8$ (light green). As is apparent from the plots of Fig. 1, the T^+_{cc} and $T^-_{\bar{c}\bar{c}}$ peaks broaden with increasing density, this effect being amplified at larger P_0. Also, notable differences arise between the two families of potentials only for small values of P_0, a feature that was discussed in Ref. [11]. The relative peak positions and widths of $T^-_{\bar{c}\bar{c}}$ and T^+_{cc} depend on both P_0 and the density, and are clearly different, due to the charge-conjugation breaking induced by the presence of the nuclear medium. For sufficiently large P_0 and density, the $T^-_{\bar{c}\bar{c}}$ maximum consistently lies at lower energy than the T^+_{cc} maximum, whereas at small P_0 and low density the

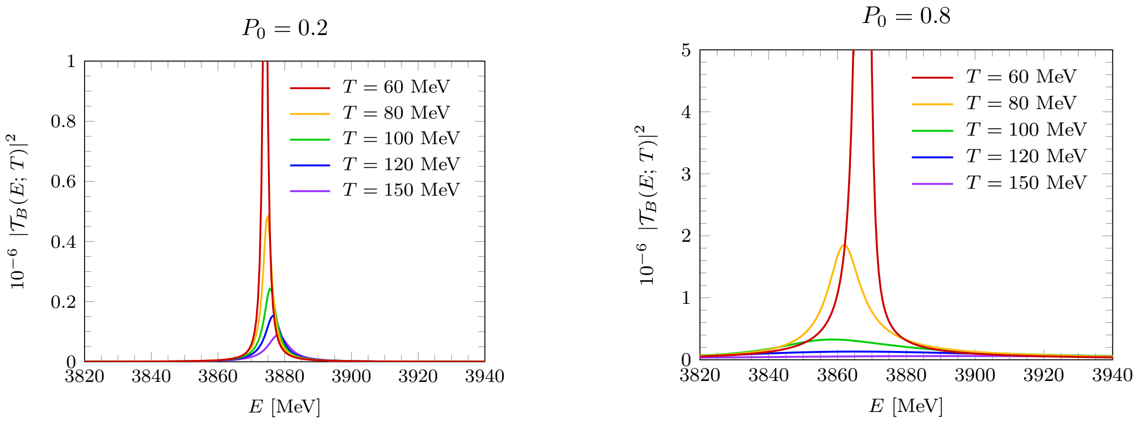

Figure 2: $\overline{D}^*\overline{D}$ (solid lines) and D^*D (dashed lines) modulus square amplitudes at finite temperature obtained by solving the on-shell BSE using the $V_B(s)$ potential, for vacuum molecular probabilities $P_0 = 0.2$ (left panel) and $P_0 = 0.8$ (right panel), and for different values of the temperature up to the crossover (different colors).

separation is barely visible. This is due to the fact that for small P_0 both $T_{cc}^{\pm}$ states couple marginally to the respective two-meson channels. Regarding the widths, $T_{\bar{c}\bar{c}}^{-}$ tends to be narrower than T_{cc}^{+}, effect that is clearer when P_0 and the density are high, although this contrast is less pronounced than the shift in peak positions and is hard to discern directly in Fig. 1. Overall, the two states respond quite differently in nuclear matter, with their nuclear-density dependent properties highly sensitive to the molecular probability in the free space.

Next, in Fig. 2 we show the temperature evolution of the $T_{cc}^{\pm}$ states (recall that the temperature modifications do not distinguish the charge-conjugated partners) for two different values of the T_{cc} compositeness (P_0), computed using the on-shell BSE with the type-B interaction kernel and the loop of Eq. (6). We do not show the results arising from V_A, since similar qualitative differences also arise in this case for small P_0, as discussed before. Concerning the temperature dependence at small and large P_0, we find that the amplitude evolves much faster with T for high molecular probabilities. For large P_0 (right plot) the width grows markedly with temperature, and the $T_{cc}(3875)^+$ melts for $T \gtrsim 100$ MeV. In contrast, for a small molecular component, $P_0 = 0.2$ (left column), the temperature effects are milder. In this case, we observe that the quasiparticle peak shifts to higher energies and only slowly broadens with T, unlike the rapid melting at large P_0. Altogether, these trends indicate that measurements of in-medium scattering amplitudes at finite temperature can provide valuable constraints on the molecular content of the $T_{cc}(3875)^+$.

In summary, we find a clear dependence of the dense nuclear medium and high temperature results on the internal structure of these exotic states. Measurements of in–medium $D^{(*)}D^{(*)}$ scattering amplitudes at finite temperature (RHIC, LHC) and at high density (CBM and PANDA at FAIR, including fixed–target $\bar{p}$–nucleus reactions), combined with comparative studies of particle vs. antiparticle channels, can tightly constrain the molecular content of the $T_{cc}(3875)^+$ and $T_{\bar{c}\bar{c}}(3875)^-$, discriminating molecular from compact configurations.

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