CHARMONIUM SPECTRAL FUNCTIONS IN $\bar{p}A$ COLLISION*

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We study the in-medium propagation of low-lying charmonium states: J/Ψ, $\Psi(3686)$, and $\Psi(3770)$ in a $\bar{p}$ Au 10 GeV collision. This energy regime will be available for the PANDA experiment. The time evolution of the spectral functions of the charmonium states is studied with a BUU-type transport model. We observe a substantial effect of the medium in the dilepton spectrum.

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1. Introduction

The observation of charmonium in vacuum and in medium is an important goal of the PANDA Collaboration at the future FAIR complex. In antiproton induced reactions, a large number of charmed states are expected to be created. Furthermore, $\bar{p}A$ reactions are best suited to observe charmed particles in nuclear matter, since the matter is in this case much simpler than the one created in a heavy-ion collision.

The spectral functions of the J/Ψ, $\Psi(3686)$, and $\Psi(3770)$ vector mesons are expected to be modified in a strongly interacting environment. In our transport model of the BUU-type, the time evolution of single-particle distribution functions of various hadrons are evaluated within the framework of a kinetic theory. The $\Psi(3770)$ meson is already a broad resonance in vacuum, while the J/Ψ, $\Psi(3686)$ mesons may acquire a noticeable width in nuclear matter [1]. Therefore, one has to propagate properly the spectral functions of these charmonium states. This is the main goal of our paper. Similar investigation has been carried out in [2].

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2. Off-shell transport of broad resonances

If we create a particle in a medium with in-medium mass through the evolution, it should regain its vacuum mass, when it reaches vacuum. If we use a local density approximation for changing its mass, we clearly break the energy conservation. We need for the propagation of off-shell particles a more sophisticated method. We can describe the in-medium properties of particles with an "off-shell transport". These equations are derived by starting from the Kadanoff–Baym equations for Green's functions of particles. Applying first-order gradient expansion after a Wigner transformation [3,4], one arrives at a transport equation for the retarded Green's function. In the medium, particles acquire a self-energy $\Sigma(x,p)$ which depends on the position and momentum as well as the local properties of the surrounding medium. Particle properties are described by their spectral function being the imaginary part of the retarded propagator

$$\mathcal{A}(p) = -2\,\mathrm{Im}\,G^{\mathrm{ret}}(x,p) = \frac{\hat{\Gamma}(x,p)}{\left(E^2 - \vec{p}^{\,2} - m_0^2 - \mathrm{Re}\,\Sigma^{\mathrm{ret}}(x,p)\right)^2 + \frac{1}{4}\hat{\Gamma}(x,p)^2}\,, \tag{1}$$

where the resonance widths Γ and $\hat{\Gamma}$ are related via $\hat{\Gamma}(x,p) = -2\,\mathrm{Im}\,\Sigma^{\mathrm{ret}} \approx 2m_0\Gamma$, and m_0 is the vacuum pole mass of the respective particle.

To solve numerically the Kadanoff–Baym equations, one may exploit the test-particle Ansatz for the retarded Green's function [3, 4]. This function can be interpreted as a product of particle number density multiplied with the spectral function $\mathcal{A}$.

The relativistic version of the equation of motion have been derived in [3]

$$\frac{\mathrm{d}\vec{x}}{\mathrm{d}t} = \frac{1}{1-C}\frac{1}{2E}\left(2\vec{p} + \vec{\partial}_p\,\mathrm{Re}\,\Sigma^{\mathrm{ret}} + \frac{m^2 - m_0^2 - \mathrm{Re}\,\Sigma^{\mathrm{ret}}}{\hat{\Gamma}}\vec{\partial}_p\hat{\Gamma}\right), \tag{2}$$

$$\frac{\mathrm{d}\vec{p}}{\mathrm{d}t} = -\frac{1}{1-C}\frac{1}{2E}\left(\vec{\partial}_x\,\mathrm{Re}\,\Sigma^{\mathrm{ret}} + \frac{m^2 - m_0^2 - \mathrm{Re}\,\Sigma^{\mathrm{ret}}}{\hat{\Gamma}}\vec{\partial}_x\hat{\Gamma}\right), \tag{3}$$

$$\frac{\mathrm{d}E}{\mathrm{d}t} = \frac{1}{1-C}\frac{1}{2E}\left(\partial_t\,\mathrm{Re}\,\Sigma^{\mathrm{ret}} + \frac{m^2 - m_0^2 - \mathrm{Re}\,\Sigma^{\mathrm{ret}}}{\hat{\Gamma}}\partial_t\hat{\Gamma}\right), \tag{4}$$

with the renormalization factor

$$C = \frac{1}{2E}\left(\partial_E\,\mathrm{Re}\,\Sigma^{\mathrm{ret}} + \frac{m_n^2 - m_0^2 - \mathrm{Re}\,\Sigma^{\mathrm{ret}}}{\hat{\Gamma}}\partial_E\hat{\Gamma}\right). \tag{5}$$

Above, $m = \sqrt{E^2 - \vec{p}^{\,2}}$ is the mass of an individual test particle. The Σ^{ret} self-energy is considered to be a function of the n baryon density, E, energy and $\vec{p}$, momentum.

Change of the m test-particle mass can be more clearly seen by combining Eqs. (3) and (4) to result in

$$\frac{\mathrm{d}m^2}{\mathrm{d}t} = \frac{1}{1-C}\left(\frac{\mathrm{d}}{\mathrm{d}t}\mathrm{Re}\,\Sigma^{\mathrm{ret}} + \frac{m^2 - m_0^2 - \mathrm{Re}\,\Sigma^{\mathrm{ret}}}{\hat{\Gamma}}\frac{\mathrm{d}}{\mathrm{d}t}\hat{\Gamma}\right), \tag{6}$$

with the comoving derivative $\mathrm{d}/\mathrm{d}t \equiv \partial_t + \vec{p}/E\vec{\partial}_x$. The vacuum spectral function is recovered when the particle leaves the medium [8].

The equations of motion of the test particles have to be supplemented by a collision term which couples the equations of the different particle species. It can be shown [4] that the collision term has the same form as in the standard BUU treatment.

In our calculations,we employ a simple form of the self-energy of a vector meson V

$$\mathrm{Re}\Sigma_V^{\mathrm{ret}} = 2m_V\Delta m_V\frac{n}{n_0}\,, \tag{7}$$

$$\mathrm{Im}\Sigma_V^{\mathrm{ret}} = -m_V\left(\Gamma_V^{\mathrm{vac}} + \Delta\Gamma_V\frac{n}{n_0}\right). \tag{8}$$

Equation (7) describes a "mass shift" $\Delta m = \sqrt{m_V^2 + \mathrm{Re}\,\Sigma_V^{\mathrm{ret}}} - m_V \approx \Delta m_V\frac{n}{n_0}$. The imaginary part contains the vacuum width Γ_V^{vac}. The second term in Eq. (8) results from the collision broadening. The parameters $(\Delta m_V, \Delta\Gamma_V)$ are taken from [1] and are the following:

Charmonium	Δm_V	$\Delta\Gamma_V$
J/Ψ	-15 MeV	20 MeV
$\Psi(3686)$	-100 MeV	20 MeV
$\Psi(3770)$	-140 MeV	20 MeV

If a meson is generated at normal density, its mass is distributed in accordance with the in-medium spectral function. If the meson propagates into a region of higher density, then the mass will be lowered according to the action of $\mathrm{Re}\,\Sigma^{\mathrm{ret}}$ in Eq. (6). However, if the meson comes near the threshold, the width $\hat{\Gamma}$ becomes very small and the second term of the right-hand side of Eq. (6) dominates, so reverses this trend leading to an increase of the mass. This method is energy conserving. We note that the propagation of ω and ρ mesons at the HADES energy range have been investigated in [5] with the same method.

3. Results

We applied Eqs. (2)–(4) to propagate the test particles for charmonium states using our BUU code [5–7]. The cross section for charmonium production in the different channels has no relevance in this study, since we do not calculate the background and we will not add absolute normalization.

In the left panel of Fig. 1, we show how the masses of test particles representing charmonium mesons are developing in $\bar{p}A$ collisions. Note that for getting a better overview in the figure, we have shifted the mass of $\Psi(3686)$ downwards by 200 MeV. At the end of the collision process, where the density is very low, the masses reach the vacuum value as it should be. The mass of $\Psi(3770)$ spreads even in the end of the collision due to the substantial vacuum width. Most of the time the masses of these mesons are either shifted downwards or at their vacuum value. The transition period between them is rather short due to the large charmonium velocity and the relatively narrow surface. The evolution of average density felt by the charmonium states is shown in the right panel of Fig. 1. This shows the same thing: the transition from the dense state to the vacuum one is approximately 4 fm/c.

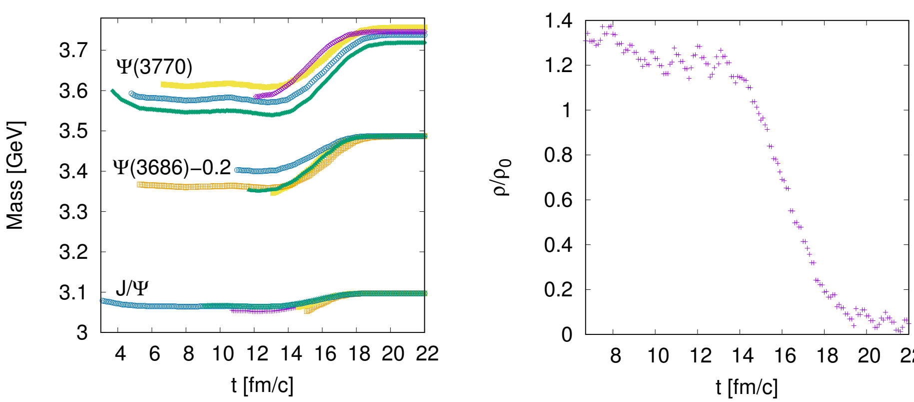

Fig. 1. In the left panel, the evolution of the masses of J/Ψ, $\Psi(3686)$, $\Psi(3770)$ is shown. For the better visibility, the masses of the $\Psi(3686)$ are shifted downwards artificially by 200 MeV. In the right panel, we show the average density as a function of time felt by the charmonium states.

In Fig. 2, we show the mass spectra of $\Psi(3770)$ at different times and also the time-integrated mass spectrum. At each time, the mass spectrum can be very well described by a Breit–Wigner form. However, the time-integrated spectrum has two peaks. One corresponds to the vacuum value, the other one to the dense phase. The valley between them is due to the short transition time.

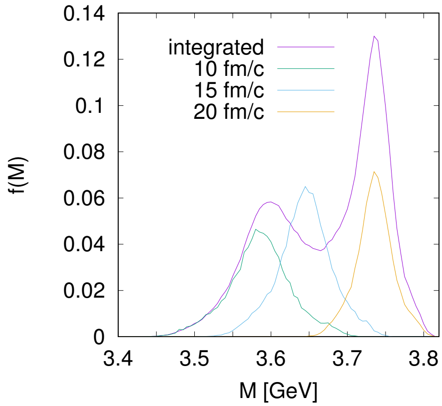

Fig. 2. We show the mass spectra of $\Psi(3770)$ at 10, 15 and 20 fm/c times, and also the time-integrated mass spectra normalized arbitrarily.

In Fig. 3, we show the charmonium contributions to the dilepton spectrum in a $\bar{p}$ Au 10 GeV collision. We compare the in-medium and the vacuum calculations. Since we did not introduce detector resolution, the vacuum calculations for J/Ψ and for $\Psi(3686)$ result in discrete lines. In the in-medium calculations, we can see the two-peak structures for each mesons, however, for the J/Ψ the effect is negligible, since its mass shift is very small. For the other two states, both peaks should be seen in experiment. In [2], this effect is not there since they do not have a mass shift for the charmonium states.

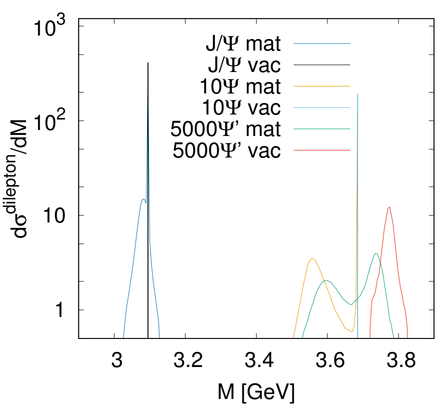

Fig. 3. Charmonium contribution to the dilepton spectra taking into account the in-medium modifications and compared it with calculations using only vacuum properties. The yields of $\Psi(3770)$ are multiplied by 5000, and the yields of $\Psi(3686)$ are multiplied by 10.

4. Summary

We calculated the charmonium contribution to the dilepton spectra. We have shown that via their dileptonic decay, there is a good chance to observe the in-medium modification of the higher charmonium states: $\Psi(3686)$, $\Psi(3770)$ in a $\bar{p}$ Au 10 GeV collision. However, we have to investigate whether the background (Drell–Yan process and the open charm decay) in the dilepton yield allows us to see this effect. This energy regime will be available by the PANDA experiment.

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