Novel mechanism for J/ψ disintegration in relativistic heavy ion collisions

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In this paper we discuss the possibility of J/ψ disintegration due to $Z(3)$ domain walls that are expected to form in a QGP medium. These domain walls give rise to a localized color electric field, which disintegrates J/ψ, on interaction, by changing its color composition and simultaneously exciting it to higher states of $c\bar{c}$ system.

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I. INTRODUCTION

The ongoing relativistic heavy ion collision experiments at the Relativistic Heavy-Ion Collider (RHIC) and the Large Hadron Collider (LHC) have provided very valuable insight in the search for quark-gluon plasma (QGP), a new phase of matter. QGP is essentially the deconfined phase of QCD, where free quarks and gluons exist in thermal equilibrium. Matsui and Satz [1] proposed that due to the presence of this medium, potential between $q\bar{q}$ is Debye screened, resulting in the swelling of quarkonia. If the Debye screening length of the medium is less than the radius of quarkonia, then $q\bar{q}$ may not form bound states. This is the conventional mechanism of quarkonia disintegration. Due to this melting, the yield of quarkonia will be *suppressed*. This was proposed as a signature of QGP and has been observed experimentally [1,2]. However, there are other factors too that can lead to the suppression of J/ψ, which has made it impossible to use J/ψ suppression as a clean signal for QGP.

In this paper, we propose a novel mechanism of quarkonia disintegration via QCD $Z(3)$ domain walls. These walls appear as topological defects due to spontaneous breaking of $Z(3)$ symmetry in QGP [3–5]. The thermal expectation value of Wilson loop (Polyakov loop) acts as the order parameter for confinement-deconfinement phase transition, taking zero value in the confining phase (corresponding to infinite free energy of a test quark) and a nonzero value in the QGP phase (with finite free energy of a test quark). Polyakov loop transforms nontrivially under $Z(3)$ transformations, and hence its nonzero expectation value leads to spontaneous breaking of $Z(3)$ symmetry in the QGP phase [6,7]. With the possibility of realization of the QGP phase in RHIC and LHC experiments, we have the real opportunity to study topological domain walls, resulting from this spontaneous $Z(3)$ symmetry breaking, in the laboratory. The formation and evolution of these walls have been discussed in the context of RHIC experiments [8,9]. The associated QGP string formation [10] has also been discussed by some of us. It is important to mention here that such topological defects invariably form during a phase transition. The formation process of topological defects is governed by formation of a sort of domain structure during a phase transition and is usually known as the *Kibble mechanism* [11]. The network of defects formed depends on the details of phase transition only through the correlation length. In fact defect distribution, e.g., defect density, per correlation volume is universal and depends only on the symmetry-breaking pattern and space dimensions.

Questions have been raised about the *reality* of these $Z(3)$ domains. The existence of these $Z(3)$ vacua becomes especially a nontrivial issue when considering the presence of dynamical quarks. The effect of quarks on $Z(3)$ symmetry and $Z(3)$ interfaces, etc., has been discussed in detail in the literature [12,13]. It has also been argued that the $Z(3)$ symmetry becomes meaningless in the presence of quarks [12]. Another viewpoint, as advocated in many papers, asserts that one can take the effect of quarks in terms of explicit breaking of $Z(3)$ symmetry [13–15]. We follow this approach and assume that the effects of dynamical quarks can be incorporated by introducing explicit symmetry-breaking terms in the effective potential for the Polyakov loop. This makes the $Z(3)$ domain wall dynamical with the pressure difference between the two different vacua being nonzero. This will lead to asymmetric profile of the Polyakov loop. In the present paper, we ignore these *asymmetry* effects due to dynamical quarks and will continue using $Z(3)$ interfaces without any explicit symmetry-breaking term. In a future work we come back to include the effects of explicit symmetry breaking.

We mention recent lattice results in Ref. [16], which indicate the strong possibility of the existence of these $Z(3)$ domains at high temperatures in the presence of dynamical quarks. These results suggest that (metastable) $Z(3)$ domains appear at temperatures above about 700 MeV. However, we stress that it does not look appropriate to take these results as conclusive, especially the quantitative part. Thus one would like to consider the possibility that $Z(3)$ vacua may persist for somewhat lower temperatures, as discussed in this paper. In any case, the mechanism discussed here provides an additional source of disintegration for J/ψ even at high temperature. It is important to note that our mechanism will lead to disintegration of Υ also, which will be relevant even at 700 MeV. We will present this study in a future work.

In the case of the early universe, these $Z(3)$ walls can lead to baryon inhomogeneity generation [17]. It was shown in Ref. [18] that background gauge field A_0 associated with *generalized* $Z(N)$ interfaces can lead to spontaneous CP violation in the Standard Model (SM), Minimal Supersymmetric Standard Model (MSSM), and Supersymmetry

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(SUSY) models, which, in turn, can lead to baryogenesis in the early universe. A detailed quantitative analysis of this spontaneous CP violation was done in Ref. [19], in the context of quark-antiquark scattering from $Z(3)$ walls in the QGP phase. The main approach followed in Refs. [18,19] was based on the assumption that the profile of the Polyakov loop order parameter $l(x)$ corresponds to a sort of condensate of the background gauge field A_0 (in accordance with the definition of the Polyakov loop). This profile of the background gauge field can be calculated from the profile of $l(x)$. Such a gauge field configuration in the Dirac equation leads to different potentials for quarks and antiquarks, leading to spontaneous CP violation in the interaction of quarks and antiquarks from the $Z(3)$ wall. This spontaneous CP violation was first discussed by Altes *et al.* [18,20] in the context of the universe and in Ref. [21] for the case of QCD. In Ref. [19], the profile of Polyakov loop $l(x)$ between different $Z(3)$ vacua was used (which was obtained by using specific effective potential for $l(x)$ as discussed in Refs. [14,15]) to obtain the profile of A_0. This background A_0 configuration acts as a potential for quarks and antiquarks. It was shown in Ref. [19] that the quarks have significantly different reflection coefficients than antiquarks and the effect is stronger for heavier quarks. For a discussion of calculation of A_0 profile, see Ref. [19].

In this paper, we discuss the effect of this spontaneous CP violation on the propagation of quarkonia in the QGP medium, in particular, the J/ψ meson. J/ψ are produced in the initial stages of relativistic heavy ion collisions. As these are heavy mesons ($m \sim 3$ GeV), they are never in equilibrium with the QGP medium produced in the present heavy ion collision experiments. However, there are finite T effects (like Debye screening, etc.) affecting its motion in a thermal bath. We ignore these effects initially and comment on them towards the end. Note that if the Debye length is larger, then the conventional mechanism of J/ψ melting does not work. As we will argue, for large Debye screening, our mechanism of J/ψ disintegration works better as any possible screening of the domain wall over the relevant length scale of J/ψ will be small. If a domain wall is present in the QGP, then a J/ψ moving through the wall will have a nontrivial interaction with it. Due to the CP-violating effect of the interface on quark scattering, c and $\bar{c}$ in J/ψ experience different color forces depending on the color of the quark and the color composition of the wall. This not only changes the color composition of $c\bar{c}$ bound state (from color singlet to color octet state) but also facilitates its transition to higher excited states (for example χ states). Color octet quarkonium states are unbound (also, the χ state has larger size than J/ψ and the Debye length), and hence they will dissociate in the QGP medium. This summarizes the basic physics of our model discussed in this paper for quarkonia disintegration due to $Z(3)$ walls.

The paper is organized in the following manner. In Sec. II we discuss the interaction of J/ψ with the background gauge field A_0 arising from the profile of $l(x)$ and discuss its color excitations. Subsequently we consider spatial excitations of J/ψ and calculate the disintegration probability. Section III discusses results, and conclusions are presented in Sec. IV.

II. INTERACTION OF J/ψ WITH A $Z(3)$ WALL

In our model, J/ψ interacts with the gauge field A_0 corresponding to the $l(x)$ profile of the $Z(3)$ wall. This allows for the possibility of color excitations of J/ψ as well as the spatial excitations of its wave function. First we discuss the possibility of color excitations of J/ψ. Subsequently, we discuss spatial excitations of J/ψ.

A. Color excitation of J/ψ

We work in the rest frame of J/ψ and consider the domain wall coming and hitting the J/ψ with a velocity v along z axis. The gauge potential and coordinates are appropriately Lorentz transformed as

$$A_0(z) \to A'_0(z') = \gamma\,[A_0(z) - vA_3(z)]\,, \tag{1a}$$

$$A_3(z) \to A'_3(z') = \gamma\,[A_3(z) - vA_0(z)]\,, \tag{1b}$$

$$z = \gamma(z' + vt'). \tag{1c}$$

We assume that there is no background vector potential, $A_i(z) = 0; i = 1,2,3$. A'_3 obtained from Eq. (1b) has only z' dependence, so it does not produce any color magnetic field. Further, using the nonrelativistic approximation of the Dirac equation one can see that the perturbation terms in the Hamiltonian [say, $H^1(A'_3)$] involving A'_3 are suppressed compared to the perturbation term [$H^1(A'_0)$] involving A'_0 at least by a factor

$$\frac{H^1(A'_3)}{H^1(A'_0)} \sim \frac{v}{c}\frac{1}{m_c r_{J/\psi}}, \tag{2}$$

where $r_{J/\psi}$ is the size of the J/ψ wave function and m_c is the charm quark mass. As we will see, the largest value of v/c we consider is 0.20–0.24 (above which transition amplitude becomes too large to trust first-order perturbation approximation). With $r_{J/\psi} \simeq 0.4$ fm, the suppression factor in Eq. (2) is of order 10%. Thus we neglect perturbation due to A'_3 and only consider perturbation due to A_0 as given by Eq. (1a). We use first-order time-dependent perturbation theory to study the excitation of J/ψ due to the background A_0 profile and consider the transition of J/ψ from initial energy eigenstate ψ_i with energy E_i to the final state ψ_j with energy E_j. The transition amplitude is given by

$$\mathcal{A}_{ij} = \delta_{ij} - i\int_{t_i}^{t_f} \langle\psi_j|\mathcal{H}_{\rm int}|\psi_i\rangle e^{i(E_j-E_i)t}dt. \tag{3}$$

We take incoming quarkonia to be a color singlet state. The interaction of the quarkonia with the wall is written as

$$\mathcal{H}_{\rm int} = V^q(z'_1) \otimes \mathbb{1}^{\bar{q}} + \mathbb{1}^q \otimes V^{\bar{q}}(z'_2) \tag{4a}$$

$$\text{with} \quad V^{q,\bar{q}}(z'_{1,2}) = gA'^{q,\bar{q}}_0(z'_{1,2}), \tag{4b}$$

where $A'^{q,\bar{q}}_0(z'_{1,2})$ is the background field configuration in the rest frame of J/ψ. z'_1 and z'_2 are the coordinates of q and $\bar{q}$ in quarkonia and g is the gauge coupling. The gauge potential A_0 is taken in the diagonal gauge as

$$A_0 = \frac{2\pi T}{g}\,(a\lambda_3 + b\lambda_8)\,, \tag{5}$$

where λ_3 and λ_8 are the Gell-Mann matrices. Under CP, $A_0 \to -A_0$, and hence $A_0^{\bar{q}} = -A_0^q$. Now, both the initial and the final states have spatial, spin, and color parts. The incoming quarkonia is a color singlet while outgoing state could be a singlet or an octet. Using Eqs. (4) and (5) and extracting only the color part of interaction, we get

$$\langle\psi_{\rm out}|\mathcal{H}_{\rm int}|\psi_{\rm singlet}\rangle = \langle\psi_{\rm out}|gA_0'^{q}(z_1')\otimes\mathbb{1}^{\bar{q}}|\psi_{\rm singlet}\rangle + \langle\psi_{\rm out}|\mathbb{1}^{q}\otimes gA_0'^{\bar{q}}(z_2')|\psi_{\rm singlet}\rangle. \quad (6)$$

The color singlet state of J/ψ is written as

$$|\psi_{\rm singlet}\rangle = \frac{1}{\sqrt{3}}\left[\begin{pmatrix}1\\0\\0\end{pmatrix}^{q}\otimes\begin{pmatrix}1\\0\\0\end{pmatrix}^{\bar{q}} + \begin{pmatrix}0\\1\\0\end{pmatrix}^{q}\otimes\begin{pmatrix}0\\1\\0\end{pmatrix}^{\bar{q}} + \begin{pmatrix}0\\0\\1\end{pmatrix}^{q}\otimes\begin{pmatrix}0\\0\\1\end{pmatrix}^{\bar{q}}\right]. \quad (7)$$

If the outgoing state is also a singlet then, each term on the right-hand side of Eq. (6) is zero due to the traceless nature of A_0. Equation (3) gives $\mathcal{A}_{ij} = 1$ for ground state ($i = j$). (One will then need to resort to second-order perturbation theory for consistency.) For higher orbital states ($i \neq j$), amplitude is identically zero. A color octet state like $|r\bar{g}\rangle$ can be written as

$$|r\bar{g}\rangle = \begin{pmatrix}1\\0\\0\end{pmatrix}^{q}\otimes\begin{pmatrix}0\\0\\1\end{pmatrix}^{\bar{q}}. \quad (8)$$

For such an outgoing state, each term on the right-hand side of Eq. (6) again vanishes identically because of the diagonal form of A_0, resulting in zero transition probability. The same argument leads to zero transition probability to all other octet states with similar color content, viz. $b\bar{g}, b\bar{r}, g\bar{r}, g\bar{b}, r\bar{b}$. There are only two states with nonzero color contribution to transition probability. They are

$$|r\bar{r} - b\bar{b}\rangle = \frac{1}{\sqrt{2}}\left[\begin{pmatrix}1\\0\\0\end{pmatrix}^{q}\otimes\begin{pmatrix}1\\0\\0\end{pmatrix}^{\bar{q}} - \begin{pmatrix}0\\1\\0\end{pmatrix}^{q}\otimes\begin{pmatrix}0\\1\\0\end{pmatrix}^{\bar{q}}\right] \quad (9)$$

and

$$|r\bar{r} + b\bar{b} - 2g\bar{g}\rangle = \frac{1}{\sqrt{6}}\left[\begin{pmatrix}1\\0\\0\end{pmatrix}^{q}\otimes\begin{pmatrix}1\\0\\0\end{pmatrix}^{\bar{q}} + \begin{pmatrix}0\\1\\0\end{pmatrix}^{q}\otimes\begin{pmatrix}0\\1\\0\end{pmatrix}^{\bar{q}} - 2\begin{pmatrix}0\\0\\1\end{pmatrix}^{q}\otimes\begin{pmatrix}0\\0\\1\end{pmatrix}^{\bar{q}}\right]. \quad (10)$$

Using Eqs. (9) and (10) in conjunction with Eqs. (5), (1), and (6), we get the color part of transition probability as

$$\langle r\bar{r} - b\bar{b}|\mathcal{H}_{\rm int}|\psi_{\rm singlet}\rangle = \frac{1}{\sqrt{6}}\left(A_0^r - A_0^b\right) \quad \text{and} \quad (11a)$$

$$\langle r\bar{r} + b\bar{b} - 2g\bar{g}|\mathcal{H}_{\rm int}|\psi_{\rm singlet}\rangle = \frac{1}{\sqrt{18}}\left(A_0^r + A_0^b - 2A_0^g\right), \quad (11b)$$

where A_0^r, A_0^b, and A_0^g are the diagonal components of the matrix $A_0'(z_1') - A_0'(z_2')$. Equations (11a) and (11b) are the effective interactions that lead to the excitations of incoming J/ψ (in the color singlet state of $c\bar{c}$) to the corresponding octet state. Due to repulsive Coulombic interaction of q and $\bar{q}$ in the octet representation, one may expect that J/ψ may disintegrate while traversing through a $Z(3)$ wall purely by color excitation. However, we will see in the next section that this is not so, and one needs to also consider spatial excitation of J/ψ due to $Z(3)$ wall.

B. Spatial excitations of J/ψ

We now consider the spatial excitations. The spatial part of the states is decided by the potential between $c\bar{c}$ in J/ψ which is taken as

$$V(|\vec{r}_1 - \vec{r}_2|) = -\frac{\alpha_s C_F}{|\vec{r}_1 - \vec{r}_2|} + C_{cnf}\sigma|\vec{r}_1 - \vec{r}_2|, \quad (12)$$

where α_s is the strong coupling constant and σ is the string tension. For J/ψ, we use charm quark mass $m_c = 1.28$ GeV, $\alpha_s = \pi/12$, and $\sigma = 0.16$ GeV^2 [22,23]. C_F is the color factor depending on the representation of the $c\bar{c}$ state. $C_F = 4/3$ for singlet state, while $C_F = -1/6$ for the octet states, showing the repulsive nature of the Coulombic part of the interaction for the octet states. C_{cnf} denotes the representation dependence of the confining part of the potential. For general sources, this factor follows Casimir scaling [24] for the string tension. For J/ψ in singlet representation, $C_{cnf} = 1$ with the value of σ used here [22,23]. It is not clear what should be the value of C_{cnf} if $c\bar{c}$ are in the octet representation. As the Coulombic part of the potential is repulsive for the octet state of $c\bar{c}$ (with $C_F = -1/6$), it is not clear if there should be a confining part of the potential at all in this case for large distances. Early lattice simulations had indicated some possibility of mildly rising potential for the confining part for $q\bar{q}$ in octet representation [25]. However, recent simulations do not show any such possibility. At large distances, the net potential between a q and $\bar{q}$ in color octet state appears to be independent of distance [26]. With the repulsive Coulombic part, this implies a very small value for C_{cnf} for the confining part. For our purpose it suffices to assume that in the octet representation, J/ψ becomes unbound, having repulsive interaction at short distances.

We have seen above that the form of A_0 in Eq. (5) only allows for transition from color singlet to two of the color octet states given in Eqs. (9) and (10). As we discussed above, $c\bar{c}$ in color octet state is unbound. Thus our task should be to consider transition from initial color singlet J/ψ to unbound state of $c\bar{c}$, say in plane waves. However, this also does not look correct as the initial J/ψ (in the color singlet state) transforms to a color octet state only as it traverses the $Z(3)$ wall [as coefficients a and b in Eq. (5) undergo spatial variations]. Thus during the early part of the passage of J/ψ through the wall, it should be dominantly in the singlet state (which is a bound state) and it will be incorrect to consider transition to unbound, plane wave states of $c\bar{c}$ at this stage. Only at later stages, when the octet component is dominant, it may be appropriate to consider repulsive potential in Eq. (12) and unbound $c\bar{c}$ states for the transition probability. This means that the perturbation term should appropriately account for the growth of octet component for the potential in Eq. (12), along with a continuing singlet component with corresponding singlet potential in Eq. (12). This clearly is a complex issue, and a proper account of appropriate potential for this type of evolution of J/ψ cannot be carried out in simple approximation scheme considered here. We make a simplifying assumption that J/ψ becomes unbound only when it transforms to the octet representation *after* its interaction with the $Z(3)$ wall. Until then it is assumed to be in the color singlet representation. Thus, in the calculations of the spatial excitation of the J/ψ state below, we use the $c\bar{c}$ potential [Eq. (12)] in the color singlet representation. The underlying physics is that incoming J/ψ is in the color singlet state, and it interacts with $Z(3)$ wall, which excites it to higher state (spatial excitation), still in color singlet potential. While traversing the $Z(3)$ wall, and undergoing this spatial excitation, the J/ψ state also transforms to color octet state. The final state, after traversing the $Z(3)$ wall, is spatially excited state in color octet representation, and our calculations give probability for this final state. This final octet state is unbound and hence such excited J/ψ disintegrates. We emphasize that at this stage our aim is to point out the new possibilities of disintegration of J/ψ with $Z(3)$ walls and this simplifying assumption should not affect our qualitative considerations and approximate estimates. We hope to give a more complete treatment in future. Thus, we continue to use the color singlet potential in Eq. (12), while considering the spatial excitation of J/ψ.

Since the potential is central, we perform coordinate transformations

$$\vec{R}_{cm} = \frac{\vec{r}_1 + \vec{r}_2}{2} \quad \text{and} \quad \vec{r} = \vec{r}_1 - \vec{r}_2, \tag{13}$$

where $\vec{r}$ is the relative coordinate between q and $\bar{q}$. $\vec{R}_{cm}$ is the center of mass of J/ψ. Using Eq. (13) with Eq. (1), we get

$$A_0^r = \gamma A_0^{11}[\gamma(z_1' + vt')] - \gamma A_0^{11}[\gamma(-z_2' + vt')]. \tag{14}$$

z_1' and z_2' are written in terms of $\vec{R}_{cm}$ and $\vec{r}$. Similar expressions can be obtained for A_0^b and A_0^g. In the above coordinates, the J/ψ wave function is $\Psi(\vec{R}_{cm})\psi(\vec{r})$. For simplicity, we assume that the center-of-mass motion remains unaffected by the external perturbation. Then $\Psi(\vec{R}_{cm})$ has the plane wave solution, while $\psi(\vec{r})$ can be written $\psi(r,\theta,\phi) = \psi(r)Y_l^m(\cos\theta,\phi)$. As J/ψ is the $l = 0$ state, we have

$$\psi_i = \psi(r)Y_0^0 \quad \text{and} \quad \psi_j = \psi_n(r)Y_l^m(\cos\theta,\phi). \tag{15}$$

The radial part, $\psi(r)$, is obtained by solving radial part of Schrödinger equation with effective potential given by

$$V(r) = -\frac{\alpha_s C_F}{r} + C_{cnf}\sigma r + \frac{l(l+1)}{2\mu r^2}, \tag{16}$$

where μ is the reduced mass. When we use Eqs. (11), (14), and (15) in Eq. (3), we get one of the terms as

$$\int_{-\infty}^{\infty} \psi_j^* A_0^r \psi_i d\vec{r}_1 d\vec{r}_2 = \int_0^{\infty}\int_{-1}^{1}\int_0^{2\pi} \psi_n^*(r)Y_l^{m*}(\cos\theta,\phi)A_0^r \times Y_0^0\psi_{100}(r)r^2 dr d(\cos\theta)d\phi. \tag{17}$$

In the above equation, we have ignored the motion of the center of mass of charmonium and have considered only the relative coordinate. Under $\cos\theta \to -\cos\theta, A_0^r \to -A_0^r$ and ψ_i does not change. So if $Y_l^m(\cos\theta,\phi) = Y_l^m(-\cos\theta,\phi)$ then the right-hand side of Eq. (17) is zero. Thus we do not get any transition to a state which is symmetric under $\cos\theta \to -\cos\theta$. This has very important significance. While the color part prohibits the transition to singlet final states, the space dependence of interaction forbids the transition to the $l = 0$ state (in color octet). Thus we see that purely color excitation of J/ψ due to A_0 field of a domain wall is not possible. The excitation is possible to the first excited state of an octet (like an "octet χ" state). As the excited state will have a radius larger than the $l = 0$ state it is more prone to melting in the medium (though with color octet composition, the final state becomes unbound anyway).

III. RESULTS

We numerically compute the integral given in Eq. (3) with various parameters given after Eq. (12). The profile of A_0 is calculated from the profile of the Polyakov loop order parameter for a $Z(3)$ domain wall at a temperature $T = 400$ MeV (as a sample value). The details of this are given in Ref. [19]. As explained there, the resulting profile is very well fitted by the functional form $p\tanh(qx + r) + s$; see Fig. 1.

We calculated the wave functions for various states of $c\bar{c}$ with the complete potential given by Eq. (16). For the

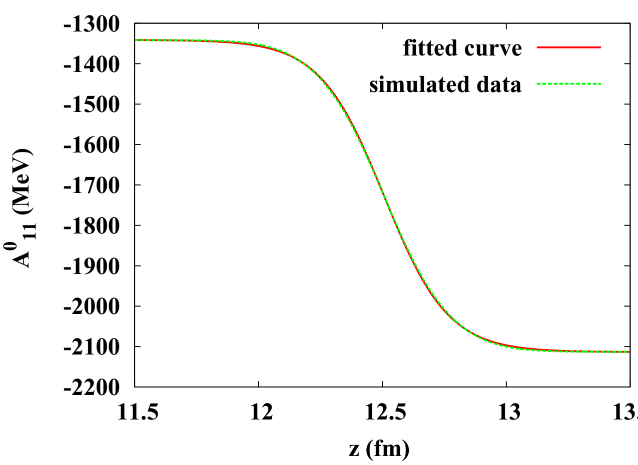

FIG. 1. (Color online) A_0 profile across the $Z(3)$ domain wall for $T = 400$ MeV. Only (1,1) component is shown. Other components are similar. See Ref. [19] for details.

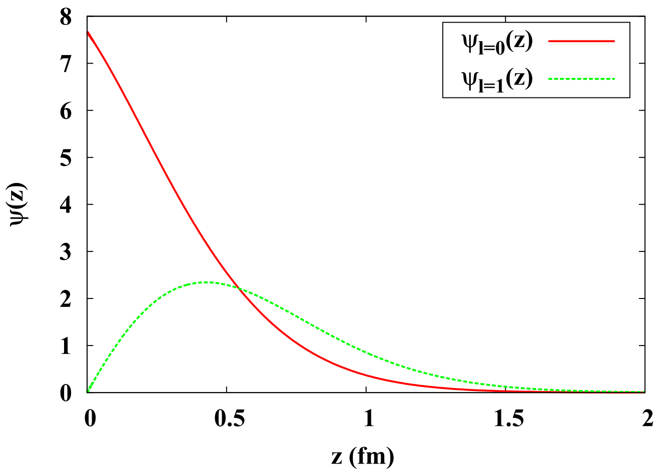

FIG. 2. (Color online) Wave functions for $J/\psi(l=0)$ and $\chi(l=1)$ states.

calculation of the wave functions for various states of $c\bar{c}$ we have used Numerov method for solving the Schrödinger equation. We have also used energy minimization technique to get the wave functions and the bound-state energy, and the results obtained by both the methods match very well. Figure 2 shows the radial part of the wave function for the $l=0,1$ states of charmonium. The bound-state contributions to the energy (excluding the rest mass of quarks) are found to be $E_0 = 0.447$ GeV for J/ψ and $E_0 = 0.803$ GeV for χ state ($l=1$). We see from Fig. 2 that the radius of J/ψ is about 0.5 fm while that for χ is about 0.8 fm. Debye length in QGP at $T = 200$ MeV is $r_d \sim 0.6$ fm and smaller at higher temperatures. Thus χ state is unstable and it should melt easily in the medium (apart from the fact that in color octet state it also becomes unbound). Figure 3 shows the combined probability of transition to both the color octet χ states [Eqs. (9) and (10)] for an incoming J/ψ with different velocities moving normal to the domain wall. As we see, the

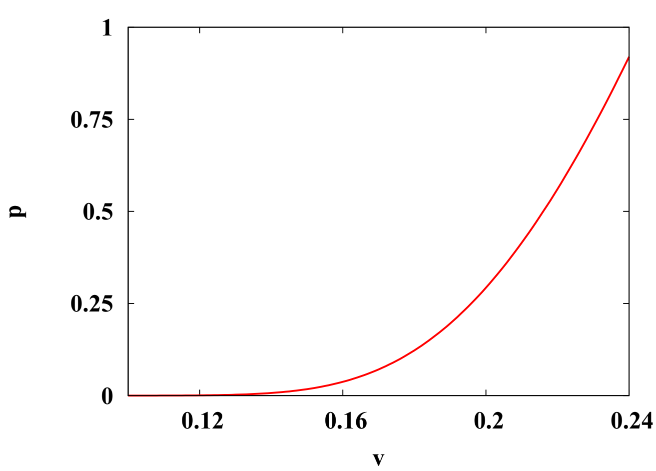

FIG. 3. (Color online) Probability p of transition of J/ψ to color octet χ states vs its velocity v. Note that the probability rapidly rises with v.

probability rapidly rises as a function of velocity. However, for large velocities the probability of transition becomes large, making first-order perturbation approximation insufficient, and one needs more reliable estimates. Thus, the plot in Fig. 3 should be trusted only for small velocities. Nonetheless, the trend at higher velocities strongly suggests that most of J/ψ will disintegrate while interacting with $Z(3)$ walls.

IV. CONCLUSIONS

These results show that on interaction with a $Z(3)$ domain wall, a J/ψ particle will make an excitation to a higher orbital state in color octet representation which is unbound and will readily melt in the surrounding QGP medium. At higher energies, the transition probability keeps increasing, making the first-order perturbation theory inapplicable and the results are not trustworthy. Nonetheless, this implies that at higher energies, almost all J/ψ are expected to disintegrate in this manner. This strong P_T dependence of J/ψ disintegration probability is a distinctive signature of our model wherein the probability of disintegration of J/ψ is enhanced with higher P_T. This can be used to distinguish this mechanism from the conventional Debey screening suppression. A very crucial point in the entire discussion is the Debye screening of the A_0 profile of the domain wall itself as it carries color. At temperature 400 MeV, the domain wall has a thickness of $\sim$1.5 fm and the Debye radius for QGP is $\sim$0.7 fm. This means that Debye screening will be effective outside a sphere of diameter $\sim$1.5 fm. So we do not expect the domain wall to be significantly Debye screened. In the above discussion, we have completely ignored the effects of a thermal bath (QGP medium) on the potential [Eq. (12)] between $c\bar{c}$ [22,27]. However, as these effects make the potential between $c\bar{c}$ weaker, the charmonium state *swells*, so it will be even easier for the interaction to break these bound states. These temperature effects will also be crucial for other heavier $q\bar{q}$ states like bottomonium as they have large binding energies. Another important aspect which has been ignored for the sake of simplicity, in the above calculations, is the question of the center-of-mass motion. This assumption is correct only in an average sense as the average force $(\Delta V/\Delta z)$ acting on c and $\bar{c}$ vanishes. This averaging is done over the thickness Δz, which is the thickness of the domain wall itself. However, as the instantaneous force $(\partial V/\partial z)$ is nonzero, there is a nonzero instantaneous acceleration of the center of mass. A more detailed analysis of the problem is required to incorporate all these details. One also needs to include the effects of dynamical quarks leading to explicit breaking of $Z(3)$ symmetry. We mention that such a disintegration of J/ψ from a color electric field may not necessarily come from a background domain wall arising in QGP medium. In a thermal medium there are always statistical fluctuations. These gluonic fluctuations will have energy of order $\sim T$. Depending on the correlation length of the fluctuation, a J/ψ passing through it may disintegrate via the mechanism discussed above. It would be interesting to study the effect of these thermal gluonic fluctuations on the spectrum of mesons.

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