J/Psi dissociation in parity-odd bubbles

Kirill Tuchin
Iowa State University, tuchin@iastate.edu

Follow this and additional works at: http://lib.dr.iastate.edu/physastro_pubs
Part of the Astrophysics and Astronomy Commons, and the Physics Commons

The complete bibliographic information for this item can be found at http://lib.dr.iastate.edu/physastro_pubs/147. For information on how to cite this item, please visit http://lib.dr.iastate.edu/howtocite.html.

This Article is brought to you for free and open access by the Physics and Astronomy at Iowa State University Digital Repository. It has been accepted for inclusion in Physics and Astronomy Publications by an authorized administrator of Iowa State University Digital Repository. For more information, please contact digirep@iastate.edu.
J/Psi dissociation in parity-odd bubbles

Abstract
We calculate the quarkonium dissociation rate in the P and CP-odd domains (bubbles) that were possibly created in heavy-ion collisions. In the presence of the magnetic field produced by the valence quarks of colliding ions, parity-odd domains generate electric field. Quarkonium dissociation is the result of quantum tunneling of quark or antiquark through the potential barrier in this electric field. The strength of the electric field in the quarkonium comoving frame depends on the quarkonium velocity with respect to the background magnetic field. We investigate momentum, electric field strength and azimuthal dependence of the dissociation rate. Azimuthal distribution of quarkonia surviving in the electromagnetic field is strongly anisotropic; the form of anisotropy depends on the relation between the electric and magnetic fields and quarkonium momentum $P_\perp$. These features can be used to explore the properties of the electromagnetic field created in heavy ion collisions.

Disciplines
Astrophysics and Astronomy | Physics

Comments

Rights
This is an open access article distributed under the Creative Commons Attribution License, which permits unrestricted use, distribution, and reproduction in any medium, provided the original work is properly cited.
We calculate the quarkonium dissociation rate in the $P$ and CP-odd domains (bubbles) that were possibly created in heavy-ion collisions. In the presence of the magnetic field produced by the valence quarks of colliding ions, parity-odd domains generate electric field. Quarkonium dissociation is the result of quantum tunneling of quark or antiquark through the potential barrier in this electric field. The strength of the electric field in the quarkonium comoving frame depends on the quarkonium velocity with respect to the background magnetic field. We investigate momentum, electric field strength and azimuthal dependence of the dissociation rate. Azimuthal distribution of quarkonia surviving in the electromagnetic field is strongly anisotropic; the form of anisotropy depends on the relation between the electric and magnetic fields and quarkonium momentum $P_{\perp}$. These features can be used to explore the properties of the electromagnetic field created in heavy ion collisions.

© 2011 Elsevier B.V. Open access under CC BY license.

1. Introduction

Solid theoretical arguments [1] and numerical calculations [2, 3] indicate a possible existence of very strong magnetic fields in heavy-ion collisions. Electromagnetic fields of such enormous intensity have never been experimentally studied and therefore present a great interest, which extends far beyond applications in the nuclear physics. What are the possible manifestations of such magnetic field? An effect that has recently attracted a lot of attention is the Chiral Magnetic Effect (CME) [1,4–7]. If a metastable $P$ and CP-odd bubble is induced by axial anomaly in the hot nuclear matter, then in the presence of external magnetic field $B_0$ the bubble generates an electric field which is parallel to the magnetic one. According to [5] the value of the electric field $E_0$ in the bubble is

$$E_0 = -N_c \sum_f \frac{e_f^2}{4\pi^2} \frac{\Theta}{N_f} B_0 = -\frac{2}{3} \frac{\alpha\Theta}{\pi} B_0$$

where the sum runs over quark flavors $f$ and we assumed that only three lightest flavors contribute. The value of the $\Theta$-angle fluctuates from event to event. CME refers to the macroscopic manifestation of this effect – separation of electric charges with respect to the reaction plane. This effect is a possible explanation of experimentally observed charge asymmetry fluctuations [8–10].

Other effects of the magnetic field that do not require existence of the parity-odd bubbles and have been recently discussed are: synchrotron radiation by fast fermions [11], polarization of the fermion spectra [11] (also known as the Sokolov–Ternov effect [12]), enhanced dilepton production [13] and azimuthal anisotropy of the quark–gluon plasma (QGP) [14,15]. These possible effects await their experimental study. Thus, strong magnetic field must have a powerful impact on the behavior of the quark–gluon plasma.

Since CME cannot be the only effect induced by the magnetic field we are motivated to look for magnetic field manifestations in other observables. We have recently pointed out in Ref. [16] that the Lorentz ionization, i.e. dissociation of a moving bound state in external magnetic field, is phenomenologically significant in relativistic heavy-ion collisions in the presence of strong [1,2] quasi-static [11] magnetic field generated by colliding ions. The ionization or, equivalently, dissociation happens because quarkonium constituents have finite probability to tunnel through the potential barrier in the presence of electric field, which appears upon boosting to the comoving frame. The dissociation rate depends on the magnetic field strength, bound state velocity and its binding energy. The dissociation rate indirectly depends on the properties of the nuclear matter by the way of dependence of the binding energy on temperature. Since at higher temperatures the binding energy is smaller, the dissociation rate is higher. Still, the Lorentz ionization can happen even if no matter is formed, provided that the magnetic field is strong enough and/or the bound state is fast enough.

E-mail address: tuchin@iastate.edu.
In this Letter we address a different type of ionization which is possible only if a metastable $P$ and CP-odd bubbles are formed and trigger emergence of electric field $E_0$ given by (1). Due to this field, there is a finite dissociation rate of quarkonium at rest in the laboratory frame. When boosted to the comoving frame quarkonium dissociation is a combined effect of the boosted electric $E_0$ and magnetic $B_0$ fields. In the comoving frame, electric and magnetic fields are directed at some angle with respect to each other depending on the quarkonium kinematics. The main goal of this Letter is to investigate the dissociation rate in this case.

2. Dissociation rate

Ionization probability of quarkonium equals its tunneling probability through the potential barrier. In the WKB approximation the later is given by the transmission coefficient and was calculated in [16]. In this method contribution of the quark spin can be easily taken into account. Another method of calculating the ionization probability, the imaginary time method [17,18], was employed in [19–21] by Popov, Karnakov and Mur. In particular, they derived in the non-relativistic approximation the pre-exponential factor that appears due to the deviation of the quark wave function from the quasi-classical approximation. Such a calculation requires matching quark wave function inside and outside the potential barrier [22]. Extension of this approach to the relativistic case is challenging due to analytical difficulties of the relativistic two-body problem. Fortunately, it was argued in [16,21] that the non-relativistic approximation provides a very good accuracy in the $\epsilon/\rho$ limit.

Magnetic field produced in heavy-ion collisions may have a complicated spatial and temporal structure (see e.g. [24]). However, we will assume that the field is constant on the scales relevant for the problem of $J/\psi$ dissociation. This approximation is supported by our recent arguments [11] that the relaxation time of the magnetic field is of the order of a few fm/c due to high electrical conductivity of the QGP. Also, spatial inhomogeneity in the transverse plane reveals itself at distances of the order of the nucleon size and perhaps even larger in view of uniformity of the matter distribution in the nuclei. A more quantitative estimate of the role of spatial and temporal dependence of the magnetic field on $J/\psi$ dissociation requires numerical solution of magneto-hydrodynamic equations and is beyond the scope of this Letter.

Given the electromagnetic field in the laboratory frame $B_0, E_0$, the electromagnetic field $B, E$ in the comoving frame moving with velocity $V$ is given by

$$E = E_0 \left( \gamma_L (b_0 + \rho_0^{-1} V \times b_0) - (\gamma_L - 1) V \frac{V \cdot b_0}{V^2} \right)$$  

$$B = B_0 \left( \gamma_L (b_0 - \rho_0 V \times b_0) - (\gamma_L - 1) V \frac{V \cdot b_0}{V^2} \right)$$

where $b_0 = B_0/B_0$ is a unit vector in the magnetic field direction, $\rho_0 = E_0/B_0 = 2\alpha(\theta)/3\pi$ (see (1)) and $\gamma_L = 1/\sqrt{1-V^2}$. It follows from (2) that

$$E = E_0 \sqrt{1 + \gamma_L^2 (b_0 \times V)^2 (1 + \rho_0^{-2})}$$  

$$B = B_0 \sqrt{1 + \gamma_L^2 (b_0 \times V)^2 (1 + \rho_0^{-2})}$$

where

$$\cos \theta = \frac{E \cdot B}{EB} = \frac{1}{\sqrt{(1 + \gamma_L^2 (b_0 \times V)^2 (1 + \rho_0^{-2})) (1 + \gamma_L^2 (b_0 \times V)^2 (1 + \rho_0^{-2}))}}$$

(4)

which we used the relativistic invariance of $E \cdot B$. It is useful to introduce dimensionless parameters $\gamma, \epsilon$ and $\rho$ as

$$\gamma = \frac{1}{\rho} \sqrt{\frac{2\rho_0}{m}}$$  

$$\epsilon = \frac{eE}{m^2} \left( \frac{m}{2\rho_0} \right)^{3/2}$$

(5)

where $m$ is quark mass and $\epsilon_0$ is quarkonium binding energy. We will treat the quarkonium binding potential in the non-relativistic approximation, which provides a very good accuracy to the dissociation rate [16,21]. The quarkonium dissociation rate in the comoving frame in the non-relativistic approximation is given by [19]

$$w = \frac{8\epsilon_0}{\epsilon} p(\gamma, \theta)C^2(\gamma, \theta)e^{-\frac{\gamma}{2(1 + \frac{\gamma}{2} \rho_0^2)}}$$

(6)

where function $g$ reads

$$g = \frac{3\tau_0}{2\gamma} \left[ 1 - \frac{1}{\gamma} \left( \frac{\tau_0^2}{\gamma^2} - 1 \right) \right]^{1/2} \sin \theta - \frac{\gamma}{3\gamma^2} \cos^2 \theta$$

(7)

and functions $P$ and $C$ are given in Appendix A. The contribution of quark spin is taken into account by replacing $\epsilon_0 \rightarrow \epsilon_0' = \epsilon_0 - \frac{\epsilon_0}{2} (V \cdot B)$ [16]. Function $g$ represents the leading quasi-classical exponent, $P$ is the pre-factor in the $S$-wave state of quarkonium and $C$ accounts for the Coulomb interaction between the valence quarks. Parameter $\tau_0$ satisfies the following equation

$$\tau_0^2 - \sin^2 \theta (\tau_0 \cosh \tau_0 - 1)^2 = \gamma^2$$

(8)

which establishes its dependence on $\theta$ and $\gamma$. Note, that in the limit $E \rightarrow 0$ the dissociation rate (6) exponentially vanishes. This is because pure magnetic field cannot force a charge to tunnel through a potential barrier.

Eq. (6) gives the quarkonium dissociation rate in a bubble with a given value of $\theta$. Its derivation assumes that the dissociation process happens entirely inside a bubble and that $\Theta$ is constant inside the bubble. Since in a relativistic heavy ion collision many bubbles can be produced with a certain distribution of $\Theta$’s (with average $\langle \Theta \rangle = 0$) more than one bubble can affect the dissociation process. This will result in a distractive interference leading to reduction of the CP-odd effect on quarkonium dissociation. However, if a typical bubble size $R_0$ is much larger than the size of quarkonium $R_J$, then the dissociation is affected by one bubble at a time independently of others, and hence the interference effect can be neglected. In this case (6) provides, upon a proper average, a reasonable estimate of quarkonium dissociation in a heavy ion collision. We can estimate the bubble size as the size of the sphaleron, which is of the order of the chromo-magnetic screening length $\sim 1/g^2 T$, whereas the quarkonium size is of the order $\alpha_s/\epsilon_0$. Consequently, at small coupling and below the zero-field dissociation temperature (i.e. when $\epsilon_0$ is not too small) $R_0$ is parametrically much larger than $R_J$. A more quantitative estimate of the sphaleron size $R_0 \approx 1.2/\alpha_sN_T \approx 0.4$ fm [22]; whereas for $J/\psi$ $R_J \approx \alpha_s/\epsilon_0 \approx 0.1–0.2$ fm. Thus, based on this estimate bubble interference can be neglected in the first approximation. However, since the ratio $R_J/R_0$ is actually not so small this effect nevertheless warrants further investigation.

To obtain the experimentally observed $J/\psi$ dissociation rate we need to average (6) over the bubbles produced in a given event and then over all events. To this end it is important to note that

$1$ Calculations in Refs. [1–3] yield very short relaxation time of magnetic field because they neglect the electromagnetic response of the quark-gluon medium, see [11] for details.
because the dissociation rate depends only on $\rho_0^2$ it is insensitive to the sign of the $E_0$ field or, in other words, it depends only on absolute value of $\Theta$ but not on its sign. Therefore, it stands to reason that although the precise distribution of $\Theta$'s is not known, (6) gives an approximate event average with parameter $\Theta$ representing a characteristic absolute value of the theta-angle.

3. Limiting cases

Before we proceed with the numerical calculations, let us consider for illustration several limiting cases. If quarkonium moves with non-relativistic velocity, then in the comoving frame electric and magnetic fields are approximately parallel $\theta \approx 0$, whereas in the ultra-relativistic case they are orthogonal $\theta \approx \pi/2$, see (4).

In the later case the electromagnetic field in the comoving frame does not depend on $E_0$ as seen in (3) and therefore the dissociation rate becomes insensitive to the CME. In our estimates we will assume that $\rho_0 < 1$ which is the relevant phenomenological situation. Indeed, it was proposed in [5] that $\rho_0 \sim \alpha \ll 1$ produces charge fluctuations with respect to the reaction plane of the magnitude consistent with experimental data.

1) $\Theta \gtrsim 0$, i.e. electric and magnetic fields are approximately parallel. This situation is realized in the following two cases. (i) Non-relativistic quarkonium velocities: $V \ll \rho_0$ or (ii) motion of quarkonium at small angle $\phi$ to the direction of the magnetic field $b_0$: $\phi \ll \rho_0/\eta V$. In both cases $E \approx E_0$ and $B \approx B_0$. This is precisely the case where the dissociation rate exhibits its strongest sensitivity to the strength of the electric field $E_0$ generated by the local parity violating QCD effects. Depending on the value of the $\gamma$ parameter defined in (5) we can distinguish the case of strong electric field $\gamma \gg 1$ and weak electric field $\gamma \ll 1$ [20]. In the former case, $g = (3/8)\gamma$, $P = (8/e)^{1/2}e^{-\gamma^2/2}$, $C = e^{2\gamma^2/12}$, Substituting into (6) the dissociation rate reads

$$w = \frac{8e_b}{e}\sqrt{g} \frac{e^{-2\gamma^2/2}}{e^{-2\gamma^2/2}} \quad \frac{16e_b^2 m^2}{e B_0} \sqrt{g} \frac{e^{-2\gamma^2/2}}{e^{-2\gamma^2/2}}, \quad \gamma \gg 1$$

In the later case, $g = P = C = 1$ and

$$w = \frac{8e_b E_0}{e} e^{-2\gamma^2/2} \frac{2e_b m^2}{16e_b^2 m^2} (2\gamma^2/12)^{1/2}, \quad \gamma \ll 1$$

where the electromagnetic field in the comoving frame equals one in the laboratory frame as was mentioned before.

\[2) \theta \sim \pi/2\], i.e. electric and magnetic fields are approximately orthogonal. This occurs for an ultra-relativistic motion of quarkonium $V \rightarrow 1$. In this case

$$B = E = B_0 |b_0| V \sqrt{1 + \rho_0^2}$$

This case was discussed in detail in our previous paper [16]. In particular for $\gamma \ll 1$ we get

$$w = \frac{8e_b m^2}{e E_0} \left(\frac{2e_b m^2}{16e_b^2 m^2}\right)^{1/2} e^{-2\gamma^2/2} \sqrt{\rho_0^2}$$

Due to (9) and (12) dependence of $w$ on $E_0$ is weak unless $\rho_0 \gg 1$.

4. J/$\psi$ dissociation rate

One of the most interesting applications of the formalism described in the previous sections is calculation of the dissociation rate of J/$\psi$ which is considered a litmus test of the quark–gluon plasma [25].

Let $z$ be the heavy ions collision axis; heavy-ion collision geometry implies that $\mathbf{b}_0 \cdot \mathbf{z} = 0$. The plane containing $z$-axis and perpendicular to the magnetic field direction is the reaction plane. We have

$$\mathbf{b}_0 \times \mathbf{V}^2 = V_z^2 + V_\perp^2 \sin^2 \phi$$

where $\phi$ is the angle between the directions of $\mathbf{b}_0$ and $\mathbf{V}_\perp$ and we denoted vector components in the $xy$-plane by the subscript $\perp$. We can express the components of the quarkonium velocity $\mathbf{V}$ in terms of the rapidity $\eta$ as $V_z = \tanh \eta$, $V_\perp = P_\perp/\gamma (M_\perp \cosh \eta)$, where $P$ and $M$ are the quarkonium momentum and mass and $M_\perp^2 = M^2 - P_\perp^2$.

Results of numerical calculations are exhibited in Figs. 1–3. In Fig. 1 we show the dissociation rate of J/$\psi$ for several values of the electric field $E_0$ induced by the Chiral Magnetic Effect. Note, that the typical size of the medium traversed by a quarkonium is a few fm. Therefore, $w \sim 0.3–0.5$ fm$^{-3}$ corresponds to complete destruction of J/$\psi$’s. This means that in the magnetic field of strength $eB_0 \sim 15m^2$ all J/$\psi$’s with $P_\perp \geq 0.5 \text{ GeV}$ are destroyed independently of the strength of $E_0$. This $P_\perp$ is lower than we estimated previously in [16] neglecting the pre-factors in the dissociation rate. Since magnetic field strength decreases towards the QGP periphery, most of J/$\psi$ surviving at later times originate from

\[2\] Note, that the limit $\gamma \gg 1$ is different in $\Theta = \pi/2$ and $\Theta < \pi/2$ cases [19].
that region. Effect of the electric field $E_0$ of the parity-odd bubble
is strongest at low $P_\perp$, which is consistent with our discussion
in the previous section. The dissociation rate at low $P_\perp$ exponen-
tially decreases with decrease of $E_0$. Probability of quarkonium
ionization by the fields below $E_0 \lesssim 0.1B_0$ (i.e. $\rho_0 \lesssim 0.1$) is
exponentially small. This is an order of magnitude higher than the
estimate $\rho_0 \sim 0$ proposed in [5].

As the plasma temperature varies, so is the binding energy of
quarkonium although the precise form of the function $\varepsilon_b(T)$ is
model-dependent. The dissociation rate picks at some $\varepsilon_b^{\text{acc}} < \varepsilon_b^{\text{vac}}$ (see
Fig. 1(b)), where $\varepsilon_b^{\text{acc}}$ is the binding energy in vacuum, indi-
cating that $J/\psi$ breaks down even before $\varepsilon_b$ drops to zero, which
is the case at $B_0 = 0$. This $\varepsilon_b^{\text{vac}}$ is a strong function of $E_0$ as can
be seen in Fig. 2. It satisfies the equation $\partial w/\partial \varepsilon_b = 0$. In the case
$\gamma \ll 1$ (10) and (12) imply that

$$\varepsilon_b^{\text{vac}} = \frac{m}{2} \left( \frac{5eE}{2m^2} \right)^{2/3}, \quad \gamma \ll 1$$

At $\gamma \gg 1$ and $\theta = \pi/2$ we employ (9) to derive the condition
$(\varepsilon_b^{\text{vac}})^2 + eB\varepsilon_b^{\text{vac}}/2m - eE^2/B = 0$. In view of (11) $E \approx B$ and we obtain

$$\varepsilon_b^{\text{vac}} = \frac{eB}{4m} \left( \frac{\sqrt{16m^2/eB} + 1}{eB} - 1 \right) \approx \sqrt{eB}, \quad \gamma \gg 1$$

where in the last step we used that $eB \ll m^2$. For a given function
$\varepsilon_b(T)$ one can convert $\varepsilon_b^{\text{vac}}$ into the dissociation temperature, which
is an important phenomenological parameter.

In the absence of the CME the dissociation probability peaks in
the direction perpendicular to the direction of magnetic field $b_0$,
and in the reaction plane. Dissociation rate vanishes in the $b_0$
direction. Indeed, for $V \cdot b_0 = 0$ (3) implies that $E = 0$. This feature
is seen in the left panel of Fig. 3. At finite $E_0$ the dissociation proba-
bility is finite in the $b_0$ direction making the azimuthal distribution
more symmetric. The shape of the azimuthal distribution strongly
depends on quarkonium velocity: while at low $V$ the strongest
dissociation is in the direction of the reaction plane, at higher $V$
the maximum shifts towards small angles around the $b_0$ direction.
Extrema of the azimuthal distribution are roots of the equation
$\partial w/\partial \phi = 0$. At $\gamma \ll 1$ it yields minimum at $\phi_0 = 0$, maximum at
$\phi_0 = \pi/2$ and another maximum that satisfies the condition (ne-
glecting the spin-dependence of $\varepsilon_b$)

$$eE_0 \sqrt{1 + \gamma^2 (V_2^2 + V_4^2 \sin^2 \phi_0) (1 + \rho_0^2)} = \frac{2m^2}{3} \left( \frac{2\varepsilon_b}{m} \right)^{3/2}$$

(16)

In order to satisfy (16) $\phi_0$ must decrease when $V$ increases and
vice versa. This features are seen in the left panel of Fig. 3.

Spectrum of quarkonia surviving in the electromagnetic field is pro-
tional to the survival probability $P = 1 - \omega t$, where $t$ is the
time spent by the quarkonium in the field. Consider $P$ as a
function of the angle $\chi$ between the quarkonium velocity and the
reaction plane $\chi = \pi/2 - \varphi$. Fourier expansion of $P$ in $\chi$
reads

$$P(\chi) = \frac{1}{2} P_0 + \sum_{n=1}^{\infty} P_n \cos(n\chi), \quad P_n = \frac{1}{\pi} \int P(\chi) \cos(n\chi) d\chi.$$  (17)

Ellipticity of the distribution is characterized by the “elliptic flow”
coefficient $v_2$ defined as

$$v_2 = \frac{P_2}{P_0} = \frac{\int P(\chi) \cos(2\chi) d\chi}{\pi \langle P \rangle}$$

$$= -\frac{1}{\pi \langle P \rangle} \int \cos(2\chi) d\chi$$

(18)

where $\langle P \rangle$ denotes average of $P$ over the azimuthal angle. These
formulas are applicable only as long as $\omega t < 1$ because otherwise
there are no surviving quarkonia. In the right panel of Fig. 3 we
show $v_2(\langle P \rangle)/t$, which is independent of $t$, as a function of $P_\perp$. As
expected, in the absence of the CME $v_2$ is negative at low $P_\perp$ and
positive at high $P_\perp$. $v_2$ changes sign at $P_\perp$ that depends on the
strength of the electric field. It decreases as $E_0$ increases until at
$E_0 \simeq B_0$ it becomes positive at all $P_\perp$. Fig. 3(b) provides the low
bound for $v_2$ because $\langle P \rangle < 1$ and $t \geq 1$ fm. We thus expect that
magnetic field strongly modifies the azimuthal distribution of the
produced $J/\psi$’s. Role of the magnetic field in generation of az-
imuthal anisotropies in heavy-ion collisions has been pointed out
before in [13,14].

5. Summary

In this Letter we studied the effect of the parallel electric and
magnetic fields on the dissociation rate of quarkonia, and parti-
cularly of $J/\psi$. Our main observation is that the CME effect on
the dissociation rate is significantly different than the effect of the
pure magnetic field if $E_0 \gtrsim 0.1B_0$ that implies an estimate of the
$\Theta$-parameter: $|\Theta|/\pi \simeq 0.1/\alpha$. This is about an order of magnitude
larger than is required for the charge separation [5]. Due to the
electric field $E_0$ of the parity-odd bubble quarkonium dissociation
rate is finite at low $P_\perp$, as indicated in Fig. 1. Fortunately, $J/\psi$’s
can be measured down to very low $P_\perp$ [26,27]. The effect of the
electric field is most clear along the direction perpendicular to the
reaction plane, because magnetic component of the Lorentz force
vanishes in this direction, see Fig. 3(a).

Azimuthal distribution of dissociation rate is strongly asym-
metric in external magnetic field. The second harmonic $v_2$ of the
azimuthal distribution of survival probability is large and negative
at low $P_\perp$, while at high $P_\perp$ it is positive; zero of $v_2$ depends on
the relation between the electric and magnetic fields. According to
the preliminary experimental data $J/\psi$’s $v_2$ is either small, about
a few per cent, or zero [28,29]. Absence of such asymmetry in the
experimental data may have two reasons. (i) Magnetic field is sig-
ificantly weaker and short-lived than suggested in [1,11], which

Fig. 2. Contour plot of the dissociation rate of $J/\psi$ as a function of $e\varepsilon_b$ and $eE_0$ at $E_0 = 15m^2$. $\phi = \pi/2$ (in the reaction plane), $\eta = 0$ (midrapidity) and $P_\perp = 0.1$ GeV. Numbers inside boxes indicate the values of $w$ in fm.
is however at odds with the charge separation observations [8–10].

(ii) Almost none of $J/\psi$'s produced in the center of QGP survive. Rather they originate from the peripheral regions. The later scenario is realized if time $t$ spent by quarkonium in the field is large because the dissociation rate only linearly increases with $t$ but exponentially decreases with the decrease of the field strength toward the QGP periphery. Finally, if the bubble interference effects due to finite bubble size discussed at the end of Section 2 are important they can significantly reduce the CP-odd effect on $J/\psi$ dissociation rate.

Abundance of possible effects associated with strong magnetic field calls for a detailed experimental investigation.

Acknowledgement

This work was supported in part by the US Department of Energy under Grant No. DE-FG02-87ER40371.

Appendix A

Here we list functions $C$ and $P$ that appear in (6):

$$P = \frac{\gamma^2}{\tau_0} \left[ \left( \frac{\tau_0 \coth \tau_0 + \sinh \tau_0 \cosh \tau_0 - 2}{\tau_0} \right) \sin^2 \theta 
+ \frac{\sinh^2 \tau_0 \cos^2 \theta}{\tau_0} \right]^{-1/2}$$

(A.1)

$$C = \exp \left[ \ln \frac{\tau_0}{2\gamma} + \int_0^{\tau_0} d\tau \left( \frac{\gamma}{\xi(\tau)} - \frac{1}{\tau_0 - \tau} \right) \right]$$

(A.2)

$$\xi(\tau) = \left\{ \frac{1}{4} \left( \frac{\tau_0^2 - \tau^2}{\tau_0^2} \right)^2 \cos^2 \theta + \frac{1}{2} \left( \frac{\cosh \tau_0 - \cos \tau}{\sinh \tau_0} \right)^2 
- \left( \frac{\sinh \tau}{\sinh \tau_0} \right)^2 \right\}^{1/2} \sin^2 \theta$$

(A.3)

**References**