Charmonium transition in electromagnetic and rotational fields

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We study heavy quarkonia in the frame of potential model in electromagnetic and rotational fields. The rotation itself cannot induce charmonium dissociation, but its coupling to the magnetic field can largely enhance or reduce the Lorentz potential and therefore affects the charmonium properties strongly. The charmonium wave function is significantly broadened in the direction of the Lorentz force, which leads to the charmonium transition from strongly interacting bound state to magnetic and rotational interaction controlled bound state. The condition for the transition seems possible to be realized in high energy nuclear collisions.

Quarkonia have long been considered as a probe [1] of the new state of matter - quark-gluon plasma (QGP) which can be created in the early stage of high energy nuclear collisions. The cold and hot nuclear matter effects on quarkonium properties and the consequence in the final state of nuclear collisions are deeply investigated in literature, see recent reviews [2–5]. The study on quarkonia in medium is recently extended to including electromagnetic fields \(E\) and \(B\), since the strongest fields in nature can be generated in nuclear collisions. For instance the peak magnitude of the magnetic field can reach \(eB \approx 70m^2_e\) in Pb+Pb collisions at the Large Hadron Collider (LHC) [6–10], where \(e\) is the electron charge and \(m_e\) the pion mass. The evolution of the heavy quark pairs \(Q\bar{Q}\) [11], the quarkonium mass and shape [12–14], and the elliptic flow in the final state [11] are largely affected by the electromagnetic fields. The experimental data at the Relativistic Heavy Ion Collider (RHIC) show that, the strongest rotational field \(\omega\) in nature can also be produced in nuclear collisions. The maximum magnitude is about \(\omega \approx 0.1m^2_e\) in non-central \(Au+Au\) collisions [15–18]. Different from the electromagnetic fields which rapidly decay in time, the angular momentum conservation during the evolution of the collision system may make a more visible rotational effect on the final state.

There are two natural questions we ask ourselves. When the electromagnetic and rotational fields are strong enough, can quarkonium states be dissociated in these external fields, like the melting in hot medium? Is it possible for the dissociated heavy quarks \(Q\) and \(\bar{Q}\) to be bound again by the electromagnetic and rotational interaction? In this paper, we will focus on the \(Q\bar{Q}\) pair transition from the bound state of strong interaction to the bound state of electromagnetic and rotational interaction. Since bottom quark is too heavy to be sizeably affected by external fields, we consider only charm quarks.

As an effective theory to study bound states of heavy quarks, the non-relativistic or relativistic potential models, based on Schrödinger or Dirac equation, have been successfully used to describe quarkonium properties in vacuum [19, 20] and medium [21–23] for many years. Taking the lattice QCD simulated heavy quark potential at finite temperature [24–26], the calculated melting temperatures in potential models can describe well the sequential charmonium suppression observed in heavy ion collisions [27–30]. In this paper we will take the Schrödinger equation to calculate the \(\bar{c}c\) bound states in external electromagnetic and rotational fields. By taking the non-relativistic limit of the Dirac equation in curved space or an effective Galilean transformation, the rotational field can be introduced into quantum mechanics via a gauge potential [31]. The corresponding electromagnetic and rotational gauge potentials can be respectively taken as \(A^\mu = (E \cdot x_c, (B \times x_c)/2)\) and \(A^\mu = ((\omega \times x_c)/2, \omega \times x_c)\).

Replacing the normal derivative \(\partial_\mu\) for free particles by the covariant derivative \(\partial_\mu\), for free particles by the covariant derivative \(D_\mu = \partial_\mu - iqA_\mu\), where \(q\) is the particle charge, one obtains the dynamical equation for particles with gauge interaction. For a pair of charm quarks \(c\) and \(\bar{c}\), the Hamiltonian of the two-body system in coordinate space reads

\[
\hat{H} = \sum_i \left( \frac{(\hat{p}_c - q_iA_c)^2}{2m_c} - q_iA_i^c \right) + V_c + \hat{H}_{ss} + \hat{H}_{sf} \tag{1}
\]

with the quark mass \(m_c\) and charge \(q_c = -q_c = q\) for electromagnetic interaction and \(q_c = q_c = m_c\) for rotational interaction. The strong interaction between \(c\) and \(\bar{c}\) which provides the quark confinement is reflected in the central potential \(V_c\), it is usually chosen as the simple Cornell form \(V_c((x_c - x_c)| = -\alpha/(x_c - x_c| + \sigma|x_c - x_c|\), which describes well the low temperature limit of the lattice simulation [24–26].

The spin interaction contains two parts, the spin-spin potential \(\hat{H}_{ss} = V_s\hat{s}_c \cdot \hat{s}_c = \beta e^{-\gamma|x_c - x_c|}\hat{s}_c \cdot \hat{s}_c\), which is also supported by the lattice QCD simulation [32], and the spin-field potential

\[
\hat{H}_{sf} = -\frac{q}{m_c}B \cdot (\hat{s}_c - \hat{s}_c) - \omega \cdot (\hat{s}_c + \hat{s}_c), \tag{2}
\]

which is from the non-relativistic limit of the corresponding Dirac equation, where \(\hat{s}_c\) is the quark spin angular momentum.

To separate the two-body problem into a center-of-mass motion and a relative motion, one usually takes the coordinate and momentum transformation from \((x_c, x_c, p_c, p_c)\) to \((r = x_c - x_c, \mathbf{R} = (x_c + x_c)/2, \mathbf{p} = (p_c - p_c)/2, \mathbf{P} = p_c + p_c)\). When considering only electromagnetic or rotational field, the pseudo-momentum \(P_{ps} = \mathbf{P} + q/2(\mathbf{B} \times \mathbf{r})\) or \(\mathbf{P}\) is conserved during the evolution of the system, the separation can be done straightforwardly [12]. When both the fields are turned on, however, the mixing between the magnetic and rotational fields makes it impossible to find a conserved
momentum and then difficult to distinctly separate the two motions, and we have to consider the perturbative solution. To this end, we define
\[ \hat{H} = \hat{H}_0 + \hat{H}', \]
(3)
with
\[
\begin{align*}
\hat{H}_0 &= \hat{H}_c + \hat{H}_r, \\
\hat{H}_c &= \frac{\hat{P}^2}{2m_c} - 2\omega \cdot \hat{L} + 4m_c(\omega \times \hat{R})^2, \\
\hat{H}_r &= \frac{\hat{P}^2}{m_c} - \frac{q}{2m_c}(\hat{P}_{ps} + 4m_c(\hat{R} \times \omega)) \cdot (\hat{B} \times \hat{r}) \\
&\quad + \frac{q^2(\hat{B} \times \hat{r})^2}{4m_c} + \frac{m_c(\omega \times \hat{r})^2}{4} - qE \cdot \hat{r} \\
&\quad - \frac{q}{m_c}B \cdot (\hat{s}_c - \hat{s}_e) - \omega \cdot (\hat{s} + \hat{l}) \\
&\quad + V_e \hat{s}_c - \hat{s}_e + V_c, \\
\hat{H}' &= -m_c(\omega \times \hat{R})^2 - \frac{m_c(\omega \times \hat{r})^2}{4},
\end{align*}
\]
(4)
where \( \hat{H}_0 \) is divided into the center-of-mass part \( \hat{H}_c \) and the relative part \( \hat{H}_r \). We have introduced in analogy to the case with only electromagnetic field [12, 33] the pseudo-momentum
\[
\hat{P}_{ps} = \hat{P} + \frac{q}{2}B \times \hat{r} + 2m_c(\omega \times \hat{R})
\]
(5)
which is a conserved quantity with \( \{\hat{P}_{ps}, \hat{H}_0\} = 0 \). The operators \( \hat{s} = \hat{s}_c + \hat{s}_e, \hat{l} = \hat{r} \times \hat{p} \) and \( \hat{L} = \hat{r} \times \hat{P}_{ps} \) are respectively the total spin, relative and center-of-mass orbital angular momenta.

For the rotation, while the harmonic oscillator potential in \( \hat{H}_r \) will finally be canceled by the same term in the perturbation \( \hat{H}' \), the coupling to the total angular momentum \( \hat{j} = \hat{s} + \hat{l} \) shifts the binding energy of the \( cc \) state, and the linear velocity \( \hat{R} \times \omega \) changes the total momentum from \( \hat{P}_{ps} \) to \( \hat{P}_{ps} + 4m_c(\hat{R} \times \omega) \) and in turn contributes to the total Lorentz potential which is linear in \( \hat{r} \). Note that, the coupling between the rotation and the orbital angular momentum \( -\omega \cdot \hat{l} \) corresponds to the Coriolis force \( \hat{F} = -\nabla(\omega \times \hat{r}) = -\omega \times \hat{p} \). To guarantee the law of causality in rotational systems, the linear velocity should be under the constraint of \( |\hat{R} \times \omega| \leq 1 \).

With the conserved pseudo-momentum \( \hat{P}_{ps} \), the total wave function corresponding to \( \hat{H}_0 \) can be factorized as
\[
\Psi_0(\hat{r}, \hat{R}, \hat{s}_c, \hat{s}_e) = \Theta_0(\hat{R}, \hat{r})\psi_0(\hat{r}, \hat{s}_c, \hat{s}_e).
\]
The center-of-mass part is the eigenstate of the Hamiltonian \( \hat{H}_c \),
\[
\Theta_0(\hat{R}, \hat{r}) = e^{i\int(\hat{P}_{ps} - q(\hat{B} \times \hat{r})/2) \cdot \hat{R} + f(\hat{R})} \]
(6)
with \( f(\hat{R}) = -2m_c \int(\hat{\omega} \times \hat{R}) \cdot d\hat{R} \), and taking into account the relation
\[
\begin{align*}
\left( \hat{p} - \frac{q(\hat{B} \times \hat{R})}{2} - \frac{m_c(\hat{\omega} \times \hat{r})}{2} \right)^n &\Psi_0 \\
= &\Theta_0 \left( \hat{p} - \frac{m_c(\hat{\omega} \times \hat{r})}{2} \right)^n \psi_0, \quad n = 1, 2
\end{align*}
\]
(7)
the relative wave function is controlled by the Hamiltonian \( \hat{H}_r \),
\[
\hat{H}_r \psi_0 = \epsilon_0 \psi_0
\]
(8)
with the binding energy \( \epsilon_0 \).

When the relative wave function \( \psi_0 \) is known, the contribution from the perturbation \( \hat{H}' \) can be obtained through the standard method in quantum mechanics. To the first order in \( \hat{H}' \), the correction to the center-of-mass energy is \(-m_c(\omega \times \hat{R})^2\), and the binding energy and relative wave function are modified as
\[
\epsilon = \epsilon_0 - \epsilon_0 \left\langle \psi_0 \left| \frac{m_c(\omega \times \hat{r})^2}{4} \right| \psi_0 \right\rangle,
\]
\[
\psi_n = \psi_0 + \sum_{m \neq n} \frac{\epsilon_{mn}}{\epsilon_{0m} - \epsilon_{0n}} \psi_{0m},
\]
(9)
where \( \epsilon_{0m} \) and \( \psi_{0m} \) are the eigenvalue and eigenstate of \( \hat{H}_0 \) with quantum number \( m \), and the matrix element is defined as \( \epsilon_{mn} = \left\langle \psi_0 m | m_c(\omega \times \hat{r})^2/4 | \psi_0 m \right\rangle \). The condition to take the above perturbative method is the small matrix element \( \epsilon_{mn} \) in comparison with \( \epsilon_0 \). We will check this condition numerically.

Before numerically solving the Schrödinger equation (8), we first analyze the dissociation of a \( cc \) bound state in electromagnetic and rotational fields and the transition from a bound state of strong interaction to a bound state of electromagnetic and rotational interaction. To have a bound state, there must be a long-range potential between the \( c \) and \( \bar{c} \). In the current case it contains four terms; the total Lorentz potential including the contribution from both the magnetic and rotational fields, the electric potential, the confinement potential, and the harmonic oscillator potential. The first three terms are all linear in the distance between the \( c \) and \( \bar{c} \), and the last term is quadratic in the distance. Since the last term is proportional to \( qB \), its contribution to the bound state should be much weaker compared with the linear terms. Without loss of generality, we neglect the electric field \( E \) in the following to simplify the analysis and calculation. In this case, the \( cc \) bound state is mainly controlled by the total Lorentz potential and the confinement potential. Inspired from the directions of the magnetic and rotational fields created in nuclear collisions [8, 16], we take the fields \( B \) and \( \omega \) in the direction of \( e_z \). Under this choice, the long-range potential can be written as
\[
V_{long} \sim \frac{qB}{2m_c} \left[ \left( \hat{P}_{ps}^2 \cos \beta + 4m_c R_{\perp} \omega \sin \alpha \right) y \\
+ \left( -\hat{P}_{ps}^2 \sin \beta + 4m_c R_{\perp} \omega \cos \alpha \right) x \right] + \sigma r,
\]
(10)
where \( \alpha = (\hat{R}_{\perp}, e_z) \) and \( \beta = (\hat{P}_{ps}^2, e_z) \) describe the directions of the coordinate \( \hat{R} \) and pseudo-momentum \( \hat{P}_{ps} \) of the \( cc \) pair in the transverse plane perpendicular to the magnetic and rotational fields. Different from the strong interaction which is isotropic, the magnetic and rotational interaction is in the transverse plane and anisotropic in this plane. The total Lorentz potential depends on the relative direction between \( \hat{R}_{\perp} \) and \( \hat{P}_{ps} \). When the magnetic force \( \hat{F}_m = -\nabla[-q/(2m_c)\hat{P}_{ps} \cdot (\hat{B} \times \hat{r})] = q/(2m_c)\hat{P}_{ps} \times \hat{B} \) and rotational force \( \hat{F}_r = -\nabla[-2q(\hat{R} \times \omega) \cdot (\hat{B} \times \hat{r})] = 2q(\hat{R} \times \hat{\omega}) \times \hat{B} \) are in the same direction, the total Lorentz potential reaches the maximum, and when the two forces are opposite to each other, it reaches the
minimum. To see clearly the maximum and minimum magnetic and rotational effects on the $c\bar{c}$ bound state, we choose for instance $\alpha = \pi/2$, $\beta = 0$ ($R_\perp = R_\perp e_y$, $P_{ps}^\perp = R_\perp e_x$) corresponding to the maximum potential and $\alpha = -\pi/2$, $\beta = 0$ ($R_\perp = -R_\perp e_y$, $P_{ps}^\perp = R_\perp e_x$) corresponding to the minimum potential,

$$V_{\text{long}}^{\text{max/min}} \sim \frac{qB}{2m_e} (P_{ps}^\perp \pm 4m_e R_\parallel) y + \sigma r. \quad (11)$$

The potential with any angels $\alpha$ and $\beta$ is in between the two limits.

We now discuss the condition to form and dissociate a $c\bar{c}$ bound state in magnetic and rotational fields. In vacuum without the external fields, the $c$ and $\bar{c}$ are bound by the Cornell potential and the bound state is isotropic. In the case with maximum Lorentz potential, the Cornell potential in the direction of the Lorentz force ($-e_y$) is more and more suppressed by the increasing rotation $\omega$. When the rotation is strong enough, the Cornell potential is canceled in this direction, the wave function of the $c\bar{c}$ state is significantly broadened, and the bound state is anisotropically dissociated. When the rotation further increases, the $c$ and $\bar{c}$ move in a negative potential well, the binding energy between them becomes negative $\epsilon_0 < 0$, and they are again bound by the magnetic and rotational interaction. In the other directions they are still bound by the Cornell potential. In the case with minimum Lorentz potential, the magnetic effect is reduced by the rotational effect, it needs a stronger magnetic field to dissociate a bound state of strong interaction and to form a bound state of magnetic and rotational interaction. It is necessary to point out that, a rotation itself cannot dissociate or form a $c\bar{c}$ bound state, its effect is through the coupling to the magnetic field.

We start to numerically solve the Schrödinger equation (8) now. Considering the spin interaction, in general case the relative wave function $\psi(r, s_c, s_{\bar{c}})$ can not be factorized as a spatial part and a spin part. We use the total spin $s$ and its projection on the $z$-axis $s_z$ as the quantum numbers to describe the four independent spin states $|s, s_z\rangle$: the spin singlet state $|S\rangle = |0, 0\rangle$ and triplet states $|T_0\rangle = |1, 0\rangle$ and $|T_\pm\rangle = |1, \pm 1\rangle$. The spin terms in the Schrödinger equation become

$$\mathbf{B} \cdot (\hat{s}_c - \hat{s}_{\bar{c}})|T_\pm\rangle = 0,$$
$$\mathbf{B} \cdot (\hat{s}_c - \hat{s}_{\bar{c}})|T_0\rangle = B|S\rangle,$$
$$\mathbf{B} \cdot (\hat{s}_c - \hat{s}_{\bar{c}})|S\rangle = B|T_0\rangle,$$
$$\omega \cdot \hat{s}|T_\pm\rangle = \pm \omega |T_\pm\rangle,$$
$$\omega \cdot \hat{s}|T_0\rangle = 0,$$
$$\omega \cdot \hat{s}|S\rangle = 0,$$
$$\hat{s}_c \cdot \hat{s}_{\bar{c}}|T_\pm\rangle = \frac{1}{4}|T_\pm\rangle,$$
$$\hat{s}_c \cdot \hat{s}_{\bar{c}}|T_0\rangle = \frac{1}{4}|T_0\rangle,$$
$$\hat{s}_c \cdot \hat{s}_{\bar{c}}|S\rangle = -\frac{3}{4}|S\rangle. \quad (12)$$

While the coupling between spin and magnetic field keeps the triplet states $|T_\pm\rangle$ as the eigenstates of the Hamiltonian, it leads to a mixing between the spin singlet state $|S\rangle$ and triplet state $|T_0\rangle$. For the coupling between spin and rotational field, it does not make any mixing among the spin states, but creates a energy gap $\sim 2\omega$ between the two triplet states $|T_\pm\rangle$. The spin-spin interaction splits the singlet and triplet states.

For a two-body system in a central potential, the relative equation can further be separated into a radial part and an angular part, the binding energy is determined by the radial equation, and the solution of the angular part is the spherical harmonic function $Y_{lm}(\theta, \phi)$. Considering the direction dependence of the electromagnetic and rotational interactions, the potential between the quark and anti-quark is no longer a central one, a usual way to solve the relative equation is to expand the wave function in terms of the complete orthogonal spherical harmonic functions,

$$\psi_T^\pm (r, s, s_z) = \sum_{lm} a_{lm}^\pm u_{lm}^\pm (r) Y_{lm}(\theta, \phi)|T_\pm\rangle,$$
$$\psi^0_{T,S} (r, s, s_z) = \sum_{lm} [a_{lm}^0 u_{lm}^0 (r) Y_{lm}(\theta, \phi)|T_0\rangle + a_{lm}^S u_{lm}^S (r) Y_{lm}(\theta, \phi)|S\rangle], \quad (13)$$

where $a_{lm}^\pm$, $a_{lm}^0$ and $a_{lm}^S$ are the probability amplitudes for spin triplet and singlet states. By substituting the expansion into the relative equation (8), we derive the wave equations controlling the radial components $u_{lm}^\pm (r)$, $u_{lm}^0 (r)$ and $u_{lm}^S (r)$.

In relativistic heavy ion collisions, the strong electromagnetic field exists only in the very early stage of the collisions [8, 16] when the high momentum quarkonia are created via hard QCD process. The low momentum quarkonium production is dominated by the regeneration mechanism in the later stage. To have a possible application of our calculation in heavy ion collisions, we take into account in the following numerical calculation the high momentum quarkonia with $P_{ps}^\perp = 3$ GeV located at the central fireball with distance $R_\perp = 2$ fm from the rotation axis.

We apply the inverse power method [19] to numerically solving the radial equations for the charmonium ground states $J/\psi$ and $\eta_c$. By fitting the experimentally observed charmonium masses in vacuum [34], we fixed the parameters in the charmonium sector of the potential model: $m_c = 1.29$ GeV, $\alpha = 0.312$, $\beta = 1.982$ GeV, $\gamma = 0.174$ (GeV)$^2$ and $\gamma = 2.06$ GeV.

We first check if the perturbative expansion is convergent. We calculated the binding energy $\epsilon_0$ and its correction $\epsilon'$ to the first order for the four ground states, the relative corrections are $|\epsilon'/\epsilon_0|_{\eta_c} = 0.95\%$, $|\epsilon'/\epsilon_0|_{J/\psi} = 0.82\%$, $|\epsilon'/\epsilon_0|_{J/\psi} = 0.79\%$ and $|\epsilon'/\epsilon_0|_{J/\psi} = 0.42\%$. This means a fast convergence of the expansion.

The charmonium binding energy and shape in magnetic and rotational fields are shown in Fig.1 in the case with maximum Lorentz potential. We take $J/\psi_0$ as an example, and the other charmonium states behavior similarly. To focus on the rotational effect, the strength of the magnetic field $eB$ is fixed to be $10m_e^2$. In the beginning at $\omega = 0$, the binding energy $\epsilon$ is already reduced to about 85% of its vacuum value $\epsilon_0$ by the magnetic effect. With increasing rotation, the potential in the direction of the total
Lorentz force is more and more suppressed, the binding energy drops down monotonously. At the dissociation rotation \( \omega_d \simeq 0.29 \ m_\pi \), the binding energy approaches to zero, and the bound state is dissociated by the magnetic and rotational fields, see the straight dashed line in Fig.1. Beyond the dissociation point, the potential in the direction of the Lorentz force becomes negative, and the dissociated heavy quarks \( c \) and \( \bar{c} \) are again bound by the magnetic and rotational force. For the charmonium shape, we consider two quantities: the fluctuation \( \langle r \rangle = \int d^3r |\psi(r)|^2 \) and the three radii \( r_i = [\int d^3r |\psi(r)|^2]^{1/2}, \ i = x, y, z \). The former indicates the central symmetry breaking by the magnetic and rotational fields, and the latter describes the charmonium size in different directions. For a central potential, the wave function is central symmetric with \( \langle r \rangle = 0 \). In magnetic and rotational fields, the wave function becomes anisotropic. The Lorentz force broadens the wave function and leads to \( \langle y \rangle < 0 \). The Lorentz force induced wave function broadening is also the reason why the radial \( r_y \) is much larger than the other two radii \( r_x \) and \( r_z \). In the beginning at \( \omega = 0 \), the radial \( r_y \) is already two times the size in vacuum \( r_0 \) due to the magnetic field. With increasing rotation, \( r_y \) goes up rapidly but \( r_x \) and \( r_z \) increase smoothly. The small difference between \( r_x \) and \( r_z \) is from the harmonic oscillator potential \( \simeq \alpha^2 B^2 r_i^2 \). Note that, while the rotation itself cannot induce charmonium dissociation and transition, its coupling to the magnetic field plays an important role. At the dissociation rotation \( \omega_d \), the contribution from the rotation to the Lorentz potential is \( 4m_\pi R_{\perp} \omega_d/(P_{ps} + 4m_\pi R_{\perp} \omega_d) = 40\% \). Very different from the charmonium melting in hot medium where at the dissociation temperature \( T_d \) the charmonium size becomes infinity [27], the size in magnetic and rotational fields is still finite, because the dissociation rotation \( \omega_d \) is in fact the transition point for the \( c\bar{c} \) to change from a strongly interacting bound state to a magnetic and rotational bound state.

We now turn to calculating the \( J/\psi_0 \) dissociation line (transition line) in the \( \omega - eB \) plane. The result is shown in Fig.2. When the magnetic field is too weak, \( eB \leq 10m_\pi^2 \), the rotation under the constraint of the speed of light cannot trigger charmonium dissociation. When increasing magnetic field, the dissociation happens firstly in the case with maximum Lorentz potential. The \( J/\psi_0 \) transverse momentum and position are taken as \( P_{ps}=(3 \text{ GeV}) \ e_x \) and \( R_{\perp}=(2 \text{ fm}) \ e_y \).

\begin{figure}[h]
\centering
\includegraphics[width=0.8\textwidth]{fig1.pdf}
\caption{\( J/\psi_0 \) binding energy \( \epsilon \) and radii \( r_x, r_y, \) and \( r_z \) scaled by their vacuum values \( r_0 \) and \( r_0 \) in the limit of maximum Lorentz potential. The magnetic field is fixed to be \( eB = 10m_\pi^2 \), and the \( J/\psi_0 \) transverse momentum and position are taken as \( P_{ps}=(3 \text{ GeV}) \ e_x \) and \( R_{\perp}=(2 \text{ fm}) \ e_y \).
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=0.8\textwidth]{fig2.pdf}
\caption{\( J/\psi_0 \) dissociation (transition) lines in \( \omega - eB \) plane in the two limits of maximum (red line) and minimum (blue line) Lorentz potential. The \( J/\psi_0 \) transverse momentum and position are taken as \( P_{ps}=(3 \text{ GeV}) \ e_x \) and \( R_{\perp}=(2 \text{ fm}) \ e_y \).
\end{figure}
and transition in the frame of potential model in electromagnetic and rotational fields. 1) The rotation enhances or reduces the Lorentz potential, depending on the charmonium position in phase space; 2) The strong interaction between a pair of charm quarks is gradually canceled by the electromagnetic and rotational interaction in the direction of the total Lorentz force; 3) When the magnetic and rotational fields are strong enough, the strong interaction controlled $c\bar{c}$ bound states are dissociated and converted to the magnetic and rotational interaction controlled bound states; 4) The condition to have the charmonium dissociation and transition looks possible to be realized in heavy ion collisions.

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