Rotating chiral fermion system in the uniform magnetic field

Ren-Hong Fang\textsuperscript{1,*}

\textsuperscript{1}Key Laboratory of Particle Physics and Particle Irradiation (MOE),
Institute of Frontier and Interdisciplinary Science,
Shandong University, Qingdao, Shandong 266237, China

We calculate the thermodynamic quantities of the uniformly rotating system of chiral fermions under the background of a uniform magnetic field. All thermodynamical quantities are expanded at $B = 0$, $\Omega = 0$, and $\mu = 0$ as threefold series. The lower orders of these series are consistent with that from the approaches of thermal field theory and Wigner function. The zero temperature limit of these quantities are also discussed.

*fangrh@sdu.edu.cn
I. INTRODUCTION

The properties of Dirac fermion system have been investigated from many aspects for a long time. For a hydrodynamical system consisting of Dirac fermions under the background of electromagnetic fields, the Wigner function is an appropriate tool, which can provide a covariant and gauge invariant formulism. For massless (or chiral) fermion system with uniform vorticity and electromagnetic fields, the charge current and the energy-momentum tensor up to the second order have been obtained from Wigner function approach, including chiral anomaly equation, chiral magnetic and vortical effects. The pair production in parallel electric and magnetic fields with finite temperature and chemical potential from Wigner function approach is also investigated recently. Without external electromagnetic fields, the energy-momentum tensor and charge current of the massless fermion system up to second order in vorticity have been obtained from thermal field theory. For a uniformly rotating massless fermion system, the analytic expressions of the charge current and the energy-momentum tensor are obtained. For the massive and massless fermion systems under the background of a uniform magnetic field, the general expansions according to fermion mass, magnetic field and chemical potential are derived by the approaches of proper-time and grand partition function.

In this article, we consider a uniformly rotating chiral fermion system in a uniform magnetic field, where we ignore the interaction among the fermions and the directions of the angular velocity and the magnetic field are chosen to be parallel. In this article we will adopt the approach of normal ordering and ensemble average to calculate the thermodynamical quantities of the system. Firstly we briefly derive the Dirac equation in a rotating frame under the background of a uniform magnetic field from the Dirac equation in curved space. Then through solving the eigenvalue equation of the Hamiltonian, we can obtain a series of Landau levels, from which one can calculate the expectation value of corresponding thermodynamical quantities for each eigenstate. From the approach of ensemble average used in, the macroscopic thermodynamical quantities can be expressed as the summation over the product of the particle number (Fermi-Dirac distribution) and the expectation value in each eigenstate. We expand all thermodynamical quantities as threefold series at \( B = 0, \Omega = 0 \) and \( \mu = 0 \), where the lower orders are consistent with that from the approaches of thermal field theory and Wigner function respectively, and to our knowledge
the general orders have not been obtained before. We also calculate all quantities in zero
temperature limit, and obtain the equality of particle/energy density and corresponding
currents along $z$-axis.

The rest of this article is organized as follows. In Sec. II we briefly derive the Dirac
equation in a uniformly rotating frame. In Sec. III the Landau levels and corresponding
eigenfunctions of a single right-handed fermion are briefly listed. In Sec. IV and V we
calculate the particle current and energy-momentum tensor of the system. In Sec. VI
the zero temperature limit of the thermodynamical quantities are discussed. This article is
summarized in Sec. VII.

Throughout this article we adopt natural units where $\hbar = c = k_B = 1$. We use the
Heaviside-Lorentz convention for electromagnetism and the chiral representation for gamma
matrixes where $\gamma^5 = \text{diag} (-1, -1, +1, +1)$, which is the same as Peskin and Schroeder [16].

II. DIRAC EQUATION IN A UNIFORMLY ROTATING FRAME

In this section we start with briefly introducing the Dirac equation in curved spacetime [17], and then we apply it to a uniformly rotating frame [18].

In curved spacetime, under the background of the electromagnetic field, the Dirac equation for a single chiral fermion is

$$i\gamma^\mu D_\mu \psi(x) = 0,$$

(1)

where the covariant derivative $D_\mu$ and gamma matrices $\gamma^\mu$ are defined as

$$D_\mu = \partial_\mu + ieA_\mu + \Gamma_\mu, \quad \gamma^\mu = \gamma^a e^a_\mu.$$

(2)

The underline in $\gamma^\mu$ is used to distinguish the spacetime-dependent gamma matrixes $\gamma^\mu$ from the constant gamma matrixes $\gamma^a$, and $\Gamma_\mu = \frac{1}{8}\omega_{\mu ab} [\gamma^a, \gamma^b]$ is the affine connection. The definitions of vierbein $e^a_\mu$, metric tensor $g_{\mu\nu}$, and spin connection $\omega_{\mu ab}$ are listed as follows,

$$e^a_\mu = \frac{\partial x^\mu}{\partial X^a}, \quad e^a_\mu = \frac{\partial X^a}{\partial x^\mu}, \quad g_{\mu\nu} = \eta_{ab} e^a_\mu e^b_\nu,$$

(3)

$$\omega_{\mu ab} = g_{\alpha\beta} e^\alpha_a (\partial_\mu e^\beta_b + \Gamma^\beta_{\mu\nu} e^\nu_b),$$

(4)

$$\Gamma^\beta_{\mu\nu} = \frac{1}{2} g^{\beta\sigma} (g_{\sigma\nu,\mu} + g_{\sigma\mu,\nu} - g_{\mu\nu,\sigma}),$$

(5)

where $\eta_{ab} = \text{diag} (+1, -1, -1, -1)$ is the metric tensor in Minkowski space, $X^a$ and $x^\mu$ are the coordinates in a local Lorentz frame and in a general frame, respectively.
In curved spacetime, the vector $J_{V}^{\mu}$, axial vector $J_{A}^{\mu}$ and symmetric energy-momentum tensor $T^{\mu\nu}$ become

$$J_{V}^{\mu} = \bar{\psi} \gamma_{\mu}^{5} \psi, \quad J_{A}^{\mu} = \bar{\psi} \gamma_{\mu} \gamma^{5} \psi,$$

$$T^{\mu\nu} = \frac{1}{4} \left( \bar{\psi} i \gamma^{\mu} D^{\nu} \psi + \bar{\psi} i \gamma^{\nu} D^{\mu} \psi + \text{H.C.} \right),$$

where $D^{\mu}, \gamma^{\mu}$ in curved spacetime have replaced $\partial^{a}, \gamma^{a}$ in flat spacetime.

Now we consider a frame $\mathcal{K}$ rotating uniformly with angular velocity $\Omega = \Omega \mathbf{e}_{z}$ relative to an inertial frame $K$. The coordinates in $\mathcal{K}$ and $K$ are denoted as $x^{\mu} = (t, x, y, z)$ and $X^{a} = (T, X, Y, Z)$ respectively, which are related to each other by following transformations,

$$
\begin{align*}
T &= t \\
X &= x \cos \Omega t - y \sin \Omega t \\
Y &= x \sin \Omega t + y \cos \Omega t \\
Z &= z
\end{align*}
$$

According to Eq. (3), the metric tensor $g_{\mu\nu}$ and its inverse are

$$g_{\mu\nu} = \begin{pmatrix}
1 - (x^2 + y^2)\Omega^2 & y\Omega & -x\Omega & 0 \\
y\Omega & -1 & 0 & 0 \\
-x\Omega & 0 & -1 & 0 \\
0 & 0 & 0 & -1
\end{pmatrix},$$

$$g^{\mu\nu} = \begin{pmatrix}
1 & y\Omega & -x\Omega & 0 \\
y\Omega & y^2\Omega^2 - 1 & -xy\Omega^2 & 0 \\
-x\Omega & -xy\Omega^2 & x^2\Omega^2 - 1 & 0 \\
0 & 0 & 0 & -1
\end{pmatrix}.$$

Keeping $g_{\mu\nu}$ unchanged, the vierbein $e_{\mu}^{a}$ still has a freedom degree of an arbitrary local Lorentz transformation. We can choose $e_{\mu}^{a}$ as

$$e_{0}^{0} = e_{1}^{1} = e_{2}^{2} = e_{3}^{3} = 1, \quad e_{1}^{0} = -y\Omega, \quad e_{2}^{0} = x\Omega,$$

with zeros for other components.

Now we consider a single chiral fermion in a uniformly rotating frame under the background of a uniform magnetic field $\mathbf{B} = B \mathbf{e}_{z}$, and we choose the gauge potential in the
inertial frame as \(A^a = (0, \mathbf{A})\) with \(\mathbf{B} = \nabla \times \mathbf{A}\). The covariant derivative \(D_\mu\) and gamma matrices \(\gamma^\mu\) become
\[
D_\mu = \left( \partial_\mu - \frac{i}{2} \Omega \Sigma_3 + \Omega (yA_x - xA_y), \partial_x - ieA_x, \partial_y - ieA_y, \partial_z - ieA_z \right),
\]
and in this case the Dirac equation for a single chiral fermion can be written as
\[
i \frac{\partial}{\partial t} \psi(x) = \left[ -i \gamma^0 \gamma \cdot (\nabla - ie \mathbf{A}) - \Omega J_z \right] \psi(x),
\]
where \(e\) is the charge of the chiral fermion, \(J_z = \frac{1}{2} \Sigma_3 - i(x\partial_y - y\partial_x)\) is the z-component of the total angular momentum \(\mathbf{J}\), and the term \(-\Omega J_z\) can be naturally explained as the coupling energy of the angular momentum \(\mathbf{J}\) and the angular velocity \(\Omega\).

III. LANDAU LEVELS FOR A SINGLE RIGHT-HANDED FERMION IN A ROTATING FRAME

In the chiral representation of gamma matrices where \(\gamma^5 = \text{diag} (-1, -1, +1, +1)\), we can divide the chiral fermion field into left-handed and right-handed fermion fields, i.e. \(\psi = (\psi_L, \psi_R)^T\). Since the equations of motion for \(\psi_L\) and \(\psi_R\) decouple, we only discuss right-handed fermion field in this article. All results can be directly generalised to the left-handed case. In the following, we set \(eB > 0\) for simplicity.

The right-handed part of Eq. (14) is
\[
i \frac{\partial}{\partial t} \psi_R(x) = H \psi_R(x),
\]
where \(H, J_{R,z} = \frac{1}{2} \sigma_3 - i(x\partial_y - y\partial_x)\) is Hamiltonian and the z-component of the total angular momentum of the right-handed fermion. In this article we shall choose the symmetric gauge for the gauge potential, i.e. \(\mathbf{A} = (-\frac{1}{2}By, \frac{1}{2}Bx, 0)\). Then the explicit form of the Hamiltonian is
\[
H = -i \sigma \cdot \nabla + \frac{1}{2} eB (y\sigma_1 - x\sigma_2) - \Omega J_{R,z}.
\]
It can be proved that, these three Hermitian operators, \(H, \hat{p}_z = -i \partial_z, J_{R,z}\), are commutative with each other, then we can construct the common eigenfuntions of them. According to
the calculations for Landau levels in Appendix A, we list the common eigenfunctions and corresponding energy as follows:

When \( m = 1/2, 3/2, 5/2, \ldots \),

\[
\psi_{\lambda nmp_z} = \sqrt{\frac{n!}{(n + m - \frac{1}{2})!}} \left( \sqrt{\frac{\varepsilon B (E + p_z + m\Omega)}{2(E + m\Omega)}} e^{-\frac{\varepsilon B}{2} \rho^2} - \frac{1}{4} L_n^{m-\frac{1}{2}} e^{i(m-\frac{1}{2})\phi} \right) e^{-\frac{\varepsilon B}{2} \rho^2 + \frac{1}{4} L_n^{m+\frac{1}{2}} e^{i(m+\frac{1}{2})\phi}} \frac{e^{-iE t + i\varepsilon B z}}{2\pi},
\]

\[
E = \begin{cases} 
\lambda \sqrt{p_z^2 + 2eBn - m\Omega}, & n > 0 \\
\rho_z - m\Omega, & n = 0 
\end{cases}.
\]

When \( m = -1/2, -3/2, -5/2, \ldots \),

\[
\psi_{\lambda nmp_z} = \sqrt{\frac{n!}{(n - m + \frac{1}{2})!}} \left( -\sqrt{\frac{\varepsilon B (E + p_z + m\Omega)}{2(E + m\Omega)}} e^{-\frac{\varepsilon B}{2} \rho^2} + \frac{1}{4} L_n^{m-\frac{1}{2}} e^{i(m+\frac{1}{2})\phi} \right) e^{-\frac{\varepsilon B}{2} \rho^2 + \frac{1}{4} L_n^{m+\frac{1}{2}} e^{i(m+\frac{1}{2})\phi}} \frac{e^{-iE t + i\varepsilon B z}}{2\pi},
\]

\[
E = \lambda \sqrt{p_z^2 + 2eB \left( n - m + \frac{1}{2} \right) - m\Omega}.
\]

where \( \lambda = \pm 1 \) represent the states with positive and negative energy respectively, and \( n = 0, 1, 2, \ldots \) represent different Landau levels. The eigenfunctions \( \psi_{\lambda nmp_z} \) are denoted by the group of good quantum numbers \( (\lambda, n, m, p_z) \), which are normalized according to

\[
\int dV \psi_{\lambda nmp_z}^\dagger \psi_{\lambda n'm'p_z'} = \delta_{\lambda\lambda'} \delta_{n'n} \delta_{m'm} \delta(p_z' - p_z).
\]

\section*{IV. PARTICLE CURRENT}

In this section we consider a right-handed fermion system under the background of a uniform magnetic field \( B = B e_z \), and the system is rotating uniformly with angular velocity \( \omega = \omega e_z \). The interaction among the fermions in this system is ignored. We assume that this rotating system is in equilibrium with a reservoir, which keeps constant temperature \( T = 1/\beta \) and constant chemical potential \( \mu \).

\subsection*{A. Ensemble average}

We will calculate the macroscopic particle current of the system at the rotation axis (i.e. at \( r = 0 \)) through ensemble average approach, in which all macroscopic thermodynamical quantities are the ensemble average of the normal ordering of the corresponding field operators.
The forms of the eigenfunctions in Eqs. (18, 20) at \( r = 0 \) or \( \rho = 0 \) are simplified to
\[
\psi_{\lambda n mp_z} = \frac{e^{-iEt + izp_z}}{2\pi} \begin{pmatrix}
\sqrt{\frac{eB(E + p_z + \Omega/2)}{2(E + \Omega/2)}} \delta_{m,1/2} \\
0
\end{pmatrix} + \frac{e^{-iEt + izp_z}}{2\pi} \begin{pmatrix}
0 \\
-\frac{i\lambda eB\sqrt{n+1}}{\sqrt{(E-\Omega/2)(E+p_z-\Omega/2)}} \delta_{m,-1/2}
\end{pmatrix},
\]
(23)
which are to be used in the following calculations of ensemble average. We find that the \( z \)-component \( m \) of the total angular momentum can only take values \( \pm 1/2 \) due to the absence of the orbital angular momentum at \( r = 0 \).

For the right-handed fermion system, the field operator of the particle current at \( r = 0 \) is
\[
J^\mu = \psi^R_\sigma \sigma^\mu \psi_R.
\]
(24)
with \( \sigma^\mu = (1, \sigma) \). From the approach of ensemble average used in [13–15], the macroscopic particle current \( J^\mu \) can be calculated from \( J^\mu \) as follows,
\[
J^\mu = \langle \cdot J^\mu \cdot \rangle = \sum_{n=1}^{\infty} \sum_{\lambda} \int_{-\infty}^{\infty} dp_z e^{\beta(E_n - \lambda\Omega/2 - \lambda\mu)} + 1 \psi^\dagger_{\lambda, n, 1/2, p_z} \sigma^\mu \psi_{\lambda, n, 1/2, p_z}
+ \sum_{n=0}^{\infty} \sum_{\lambda} \int_{-\infty}^{\infty} dp_z e^{\beta(E_{n+1} + \lambda\Omega/2 - \lambda\mu)} + 1 \psi^\dagger_{\lambda, n, -1/2, p_z} \sigma^\mu \psi_{\lambda, n, -1/2, p_z}
+ \sum_{\lambda} \int_{-\infty}^{\infty} dp_z \theta(\lambda p_z) e^{\beta(|p_z| - \lambda\Omega/2 - \lambda\mu)} + 1 \psi^\dagger_{\lambda, 0, 1/2, p_z} \sigma^\mu \psi_{\lambda, 0, 1/2, p_z},
\]
(25)
where \( \langle \cdot \cdot \cdot \rangle \) means normal ordering and ensemble average of corresponding field operator \[12, 19\], \( \theta(x) \) is the step function, and we have defined \( E_n = \sqrt{p_z^2 + 2e\lambda E_n} \). The second, third, and fourth lines of Eq. (25) represent the contributions of high Landau levels with \( m = 1/2 \), all Landau levels with \( m = -1/2 \), and the lowest Landau level with \( m = 1/2 \), respectively.

We can see that the macroscopic particle current \( J^\mu \) consists of the summation over the product of the particle number (Fermi-Dirac distribution) and the expectation value in each mode described by the quantum numbers \( (\lambda, n, m, p_z) \).

B. Particle number density

Firstly we calculate the particle number density \( \rho \equiv J^0 \) of the system. Making use of
\[
\psi^\dagger_{\lambda n mp_z} \psi_{\lambda n mp_z} = \begin{cases}
\frac{eB E + \Omega/2 + p_z}{4\pi E + \Omega/2}, & m = \frac{1}{2} \\
\frac{eB E - \Omega/2 - p_z}{4\pi E - \Omega/2}, & m = -\frac{1}{2}
\end{cases}
\]
(26)
and from Eq. (25) one can obtain

\[ \rho \beta^3 = \frac{b \omega}{16 \pi^2} + \frac{1}{2} \sum_{s=\pm 1} \frac{\partial}{\partial a} g \left( a + \frac{1}{2} s \omega, b \right), \]  

(27)

where we have defined three dimensionless quantities \( a = \beta \mu \), \( b = 2 e B \beta^2 \), \( \omega = \beta \Omega \), and have defined \( g(x, b) \) as

\[ g(x, b) = \frac{b}{4 \pi^2} \int_0^\infty dy \sum_{n=0}^\infty \sum_{s=\pm 1} \left( 1 - \frac{1}{2} \delta_{n, 0} \right) \ln \left( 1 + e^{sx-\sqrt{nb+y^2}} \right). \]  

(28)

In a recent article [11], making use of Abel-Plana formula, the authors obtained the asymptotic expansion of \( g(x, b) \) at \( b = 0 \) as follows

\[ g(x, b) = \left( \frac{7 \pi^2}{360} + \frac{x^2}{12} + \frac{x^4}{24 \pi^2} \right) - \frac{b^2 \ln b^2}{384 \pi^2} - \frac{b^2}{96 \pi^2} \ln \left( \frac{e}{2 G^6} \right) - \frac{1}{2 \pi^2} \sum_{n=0}^\infty \frac{(4n+1)!!}{(4n+4)!!} B_{2n+2} C_{2n+1}(x) b^{2n+2}, \]  

(29)

where \( G = 1.28242... \) is the Glaisher number, \( B_n \) are Bernoulli numbers, and \( C_{2n+1}(x) \) is defined and expanded at \( x = 0 \) in the following,

\[ C_{2n+1}(x) = \delta_{n, 0} + \frac{1}{(4n+1)!} \int_0^\infty dy \ln y \frac{d^{4n+1}}{dy^{4n+1}} \left( \frac{1}{e^{y+x} + 1} + \frac{1}{e^{y-x} + 1} \right) = \left( \ln 4 + \gamma - 1 \right) \delta_{n, 0} + \frac{2}{(4n+1)!} \sum_{k=0}^\infty \left( 2^{4n+2k+1} - 1 \right) \zeta'(-4n - 2k) \frac{x^{2k}}{(2k)!}. \]  

(30)

Plugging Eqs. (29, 30) into Eq. (27), one can get the threefold series expansion of the particle number density at \( a = 0, b = 0, \omega = 0 \) or \( \mu = 0, B = 0, \Omega = 0 \) as follows,

\[ \rho \beta^3 = \frac{a}{6} + \frac{a^3}{6 \pi^2} + \frac{a \omega^2}{8 \pi^2} + \frac{b \omega}{16 \pi^2} - \frac{1}{\pi^2} \sum_{n=0}^\infty \frac{B_{2n+2} b^{2n+2}}{(4n+4)!!(4n)!!} \sum_{j=0}^\infty \frac{\omega^{2j}}{(2j)!2^{2j}} \times \sum_{k=0}^\infty \left( 2^{4n+2k+2j+3} - 1 \right) \zeta'(-4n - 2k - 2j - 2) \frac{a^{2k+1}}{(2k+1)!}. \]  

(31)

The lower orders \( O(b^2, \omega^2, b\omega) \) in Eq. (31) are consistent with the perturbative results in [4, 6, 7], where the authors used the approaches of thermal field theory and Wigner function respectively.
C. Particle current along \( z \)-axis

Next we calculate the space components of the particle current \( J^\mu \). According to the rotation symmetry along \( z \)-axis of the system, the \( x \) - and \( y \)-components of \( J^\mu \) vanish. The unique nonzero component is \( J^z \). Making use of

\[
\psi^\dagger_{\lambda nmp} \sigma_3 \psi_{\lambda nmp} = \begin{cases} 
\frac{eB}{4\pi^2} \frac{E+\Omega/2+p_z}{E+\Omega/2}, & m = \frac{1}{2} \\
-\frac{eB}{4\pi^2} \frac{E-\Omega/2-p_z}{E-\Omega/2}, & m = -\frac{1}{2} 
\end{cases} \tag{32}
\]

and from Eq. (25) one can obtain

\[
J^z \beta^3 = \frac{ab}{8\pi^2} + \frac{1}{2} \sum_{s=\pm 1} s \frac{\partial}{\partial a} g \left( a + \frac{1}{2} s \omega, b \right), \tag{33}
\]

which can be expanded as the threefold series at \( a = 0, b = 0, \omega = 0 \) or \( \mu = 0, B = 0, \Omega = 0 \) as follows,

\[
J^z \beta^3 = \frac{ab}{8\pi^2} + \frac{\omega^3}{12} + \frac{\omega a^2}{48\pi^2} + \frac{\omega^2}{4\pi^2} - \frac{1}{\pi^2} \sum_{n=0}^{\infty} \frac{\beta_{2n+2} b^{2n+2}}{(4n+4)!!(4n)!!} \sum_{j=0}^{\infty} \frac{\omega^{2j+1}}{(2j+1)!2^{2j+1}} \\
\times \sum_{k=0}^{\infty} \left( 2^{4n+2k+2j+3} - 1 \right) \zeta'(-4n-2k-2j-2) \frac{a^{2k}}{(2k)!}. \tag{34}
\]

When \( \omega = 0 \) or \( \Omega = 0 \) in Eq. (34), one can obtain \( J^z \beta^3 = \frac{ab}{8\pi^2} \), which is the chiral magnetic effect \cite{20,24}; When \( b = 0 \) or \( B = 0 \) and keeping the leading order of \( \omega \) in Eq. (34), one can obtain \( J^z \beta^3 = \frac{\omega}{12} \left( 1 + \frac{3a^2}{\pi^2} \right) \), which is the chiral vortical effect \cite{25,30}.

V. ENERGY-MOMENTUM TENSOR

In this section, we will calculate the energy-momentum tensor \( T^{\mu\nu} \) (at \( r = 0 \)) of the right-handed fermion system as described in Sec. [IV]. According to the rotation symmetry along \( z \)-axis, the energy-momentum tensor at \( r = 0 \) are unchanged under the rotation along \( z \)-axis, which leads to following constraints on \( T^{\mu\nu} \):

\[
T^{01} = T^{02} = T^{12} = T^{13} = T^{23} = 0, \quad T^{11} = T^{22}. \tag{35}
\]

The possible nonzero components of \( T^{\mu\nu} \) are \( T^{00} \), \( T^{11} = T^{22} \), \( T^{33} \), and \( T^{03} \).
For the right-handed fermion system, the field operator of the symmetric energy-momentum tensor at \( r = 0 \) is

\[
T_{\mu \nu} = \frac{1}{4} \left( \psi_{R}^\dagger i \sigma^\mu D^\nu_R \psi_R + \psi_{R}^\dagger i \sigma^\nu D^\mu_R \psi_R + \text{H.C.} \right),
\]  

(36)

with \( \sigma^\mu = (1, \sigma) \) and the right-handed covariant derivative \( D^\mu_R \) defined as

\[
D^\mu_R = \left( \partial_t - \frac{i}{2} \Omega_3, -\partial_x, -\partial_y, -\partial_z \right).
\]  

(37)

The macroscopic energy-momentum tensor \( \mathcal{T}_{\mu \nu} \) can be calculated from \( T_{\mu \nu} \) as follows,

\[
\mathcal{T}_{\mu \nu} = \langle : T_{\mu \nu} : \rangle
\]

\[
= \frac{1}{4} \sum_{n=1}^{\infty} \sum_{\lambda} \int_{-\infty}^{\infty} dp_z \frac{\lambda}{\epsilon^\beta(E_n - \lambda \Omega/2 - \lambda \mu)} + \frac{1}{4} \psi_{\lambda,n,1/2,p_z}^\dagger (i \sigma^\mu D^\nu_R + i \sigma^\nu D^\mu_R) \psi_{\lambda,n,1/2,p_z} + \text{H.C.}
\]  

(38)

A. Energy density

Firstly we calculate the energy density \( \varepsilon \equiv T^{00} \) of the system. Making use of

\[
\psi_{\lambda n m p_z}^\dagger \left( i \partial_t + \frac{1}{2} \Omega_3 \right) \psi_{\lambda n m p_z} = \begin{cases} \frac{eB}{8\pi^2} (E + p_z + \Omega/2), & m = \frac{1}{2} \\ \frac{eB}{8\pi^2} (E - p_z - \Omega/2), & m = -\frac{1}{2} \end{cases}
\]  

(39)

and from Eq. (25) one can obtain

\[
\varepsilon \beta^4 = \frac{ab \omega}{16\pi^2} + \sum_{s=\pm 1} \left( \frac{3}{2} - \frac{b}{\partial \theta} \right) g \left( a + \frac{1}{2} s \omega, b \right),
\]  

(40)
which can be expanded as the threefold series at \(a = 0\), \(b = 0\), \(\omega = 0\) or \(\mu = 0\), \(B = 0\), \(\Omega = 0\) as follows,

\[
\varepsilon^4 = \frac{7\pi^2}{120} + \frac{a^2}{4} + \frac{\omega^2}{16} + \frac{a^4}{8\pi^2} + \frac{3a^2\omega^2}{16\pi^2} + \frac{\omega^4}{128\pi^2} + \frac{ab\omega}{16\pi^2} + \frac{b^2\ln b^2}{384\pi^2} + \frac{b^2}{96\pi^2} \ln \left( \frac{2e^{\gamma+1}}{G^6} \right)
\]

\[
+ \frac{1}{\pi^2} \sum_{n=0}^{\infty} \frac{(4n+1)B_{2n+2}b^{2n+2}}{(4n+4)!!(4n)!!} \sum_{j=0}^{\infty} \frac{\omega^{2j}}{(2j)!2^j}
\]

\[
\times \sum_{k=0}^{\infty} (2^{4n+2k+2j+1} - 1) \zeta'(-4n-2k-2j) \frac{a^{2k}}{(2k)!}.
\]

where the logarithmic term \(b^2\ln b^2\) has been discussed in detail in [11], and its coefficient is independent of \(\omega\) in this work. It is worth noting that there would be no such logarithmic term if the un-normal ordering description of field operators was adopted [4, 31].

### B. Pressure

The pressure \(P\) of the system is \(T^{33}\). Making use of

\[
\psi_{\lambda mnp}^- (\partial_\zeta) \psi_{\lambda mnp} = \begin{cases} \frac{\varepsilon B (E+p_\gamma+\Omega/2) p_z}{8\pi^2 E+\Omega/2}, & m = \frac{1}{2} \\ -\frac{\varepsilon B (E-p_\gamma-\Omega/2) p_z}{8\pi^2 E-\Omega/2}, & m = -\frac{1}{2} \end{cases}
\]

and from Eq. (25) one can obtain

\[
P^{\beta^4} = \frac{ab\omega}{16\pi^2} + \frac{1}{2} \sum_{s=\pm 1} g \left( a + \frac{1}{2} s\omega, b \right),
\]

which can be expanded as the threefold series at \(a = 0\), \(b = 0\), \(\omega = 0\) or \(\mu = 0\), \(B = 0\), \(\Omega = 0\) as follows,

\[
P^{\beta^4} = \frac{7\pi^2}{360} + \frac{a^2}{12} + \frac{\omega^2}{48} + \frac{a^4}{24\pi^2} + \frac{a^2\omega^2}{16\pi^2} + \frac{\omega^4}{384\pi^2} + \frac{ab\omega}{16\pi^2} - \frac{b^2\ln b^2}{384\pi^2} - \frac{b^2}{96\pi^2} \ln \left( \frac{2e^{\gamma}}{G^6} \right)
\]

\[
- \frac{1}{\pi^2} \sum_{n=0}^{\infty} \frac{B_{2n+2}b^{2n+2}}{(4n+4)!!(4n)!!} \sum_{j=0}^{\infty} \frac{\omega^{2j}}{(2j)!2^j}
\]

\[
\times \sum_{k=0}^{\infty} (2^{4n+2k+2j+1} - 1) \zeta'(-4n-2k-2j) \frac{a^{2k}}{(2k)!}.
\]

One can obtain \(T^{11}\) from the traceless condition for energy-momentum tensor, \(T^{00} = 2T^{11} + T^{33}\).
C. Energy current

The energy current along $z$-axis is $\mathcal{T}^{03}$. Making use of

$$
\psi_{\lambda nmpz}^\dagger \left( -i \partial_z + \sigma_3 i \partial_t + \frac{1}{2} \Omega \right) \psi_{\lambda nmpz} = \begin{cases} 
\frac{eB}{8\pi^2} \frac{(E+p_z+\Omega/2)^2}{E+\Omega/2}, & m = \frac{1}{2} \\
-\frac{eB}{8\pi^2} \frac{(E-p_z-\Omega/2)^2}{E-\Omega/2}, & m = -\frac{1}{2}
\end{cases}
$$

and from Eq. (25) one can obtain

$$
\mathcal{T}^{03} \beta^4 = \frac{b}{8\pi^2} \left( \frac{\pi^2}{6} + \frac{\omega^2}{8} + \frac{a^2}{2} \right) + \sum_{s=\pm 1} s \left( 1 - \frac{b}{2} \frac{\partial}{\partial b} \right) g \left( a + \frac{1}{2} s \omega, b \right),
$$

which can be expanded as the threefold series at $a = 0$, $b = 0$, $\omega = 0$ or $\mu = 0$, $B = 0$, $\Omega = 0$ as follows,

$$
\mathcal{T}^{03} \beta^4 = \frac{b}{8\pi^2} \left( \frac{\pi^2}{6} + \frac{\omega^2}{8} + \frac{a^2}{2} \right) + \left( \frac{a \omega}{6} + \frac{a^3 \omega}{24 \pi^2} \right) + \sum_{n=1}^\infty \frac{n \beta_{2n+2} b^{2n+2}}{(4n+4)! (4n)!} \sum_{j=0}^\infty \frac{\omega^{2j+1}}{(2j+1)! 2^{2j+1}}
$$

$$
\times \sum_{k=0}^\infty \left( 2^{4n+2k+2j+3} - 1 \right) \zeta'(-4n-2k-2j-2) \frac{a^{2k+1}}{(2k+1)!}.
$$

Up to now, we have obtained all thermodynamical quantities of the right-handed fermion system. For left-handed fermion system, one can derive corresponding quantities from the right-handed case through space inversion: $\rho_R \rightarrow \rho_L$, $\mathcal{J}_R^z \rightarrow -\mathcal{J}_L^z$, $\varepsilon_R \rightarrow \varepsilon_L$, $P_R \rightarrow P_L$, $\mathcal{T}^{03}_R \rightarrow -\mathcal{T}^{03}_L$, $\mu_R \rightarrow \mu_L$, $B \rightarrow B$, $\Omega \rightarrow \Omega$, where the subscripts $R, L$ are used to distinguish the quantities in right-handed case from that in left-handed case.

VI. ZERO TEMPERATURE LIMIT

Now we turn to the thermodynamics of the system at zero temperature limit. When the temperature tends to be zero, with chemical potential $\mu$, magnetic field $B$, and angular velocity $\Omega$ fixed, then the three dimensionless quantities $a = \beta \mu$, $b = 2eB\beta^2$, $\omega = \beta \Omega$ all tend to be infinity. The asymptotic behavior of $g(x, b)$ as $x \rightarrow \infty$ and $b \rightarrow \infty$ has been obtained in [11],

$$
\lim_{x, b \rightarrow \infty} g(x, b) = \frac{x^2 b}{16\pi^2},
$$
From Eqs. (27, 33), one can derive the expressions of the particle density $\rho$ and the current $J^z$ at zero temperature limit as follows,

$$\rho = J^z = \frac{eB}{4\pi^2} \left( \mu + \frac{\Omega}{2} \right).$$ (49)

At zero temperature limit, due to the coupling of the spin with the magnetic field and the angular velocity, the spin alignment of all particles and antiparticles will be along $z$-axis of the system. Since these particles are right-handed, they will move along $z$-axis with the speed of light $c$ ($c = 1$ in natural unit), so it is reasonable that the particle density $\rho$ equals to the $z$-component current $J^z$ at zero temperature limit.

From Eqs. (40, 43, 46), the expressions of energy density $\varepsilon$, pressure $P$ and energy current $T^{03}$ at zero temperature limit are

$$\varepsilon = P = T^{03} = \frac{eB}{8\pi^2} \left( \mu + \frac{\Omega}{2} \right)^2.$$ (50)

The movements of the particles and antiparticles with the speed of light along $z$-axis leads to the equality of the energy density $\varepsilon$ and the energy current $T^{03}$. Since there is no energy current along the direction of the $x$- and $y$-axis, then $T^{11}$ and $T^{22}$ vanish in this system, which results in the equality of the energy density $\varepsilon$ and the pressure $P$.

VII. SUMMARY

In this article, we have investigated the thermodynamics of the uniformly rotating right-handed fermion system under the background of a uniform magnetic field through the approach of normal ordering and ensemble average, where all thermodynamical quantities are expanded as threefold series at $B = 0$, $\Omega = 0$ and $\mu = 0$. For these threefold series, our results at lower orders are consistent with that from different methods by other authors, and to our knowledge no literature has obtained the general orders. We also calculate all quantities in zero temperature limit, and obtain the equality of particle/energy density and corresponding currents along $z$-axis. Since for the chiral fermion the right-handed part decouples from the left-handed part, in this article we only considered the case of the right-handed fermion system, which can be directly generalized to the left-handed case through space inversion.
VIII. ACKNOWLEDGMENTS

We thank De-Fu Hou for helpful discussion. This work is supported by the National Natural Science Foundation of China under Grant No. 11890713.

Appendix A: Landau levels for a single right-handed fermion

The Hamiltonian for a right-handed fermion under the background of the uniform magnetic field \( B = B_0 e_z \) is

\[
H = -i \sigma \cdot (\nabla - ieA) = -i \sigma \cdot \nabla + \frac{1}{2} eB(y\sigma_1 - x\sigma_2),
\]

where we have chosen \( A = (-\frac{1}{2}By, \frac{1}{2}Bx, 0) \) for the gauge potential. One can refer to \([19, 31, 32]\) for other choices of the gauge potential.

In the following, we will solve the eigenvalue equation of \( H \) in cylindrical coordinates,

\[
H\psi = E\psi.
\]

We can see that the three Hermitian operators, \( H, \hat{p}_z = -i\partial_z, \hat{J}_z = \frac{1}{2}\sigma_3 + (x\hat{p}_y - y\hat{p}_x) \) are commutative with each other, so the eigenfunction \( \psi \) can be chosen as

\[
\psi = \begin{pmatrix} f(r) e^{i(m-\frac{1}{2})\phi} \\ g(r) e^{i(m+\frac{1}{2})\phi} \end{pmatrix} e^{izp_z},
\]

where \(-\infty < p_z < \infty \) and \( m = \pm 1/2, \pm 3/2, \pm 5/2, \ldots \) are the eigenvalues of \( \hat{p}_z \) and \( \hat{J}_z \) respectively. The explicit form of the Hamiltonian \( H \) in cylindrical coordinates is

\[
H = \begin{pmatrix} -i \frac{\partial}{\partial z} & e^{-i\phi} \left(-i \frac{\partial}{\partial r} + \frac{1}{r} \frac{\partial}{\partial \phi} - \frac{i}{2} eB r \right) \\ e^{i\phi} \left(-i \frac{\partial}{\partial r} + \frac{1}{r} \frac{\partial}{\partial \phi} + \frac{i}{2} eB r \right) & i \frac{\partial}{\partial z} \end{pmatrix},
\]

then from Eq. (A2) we can obtain two differential equations for \( f(r), g(r) \) as follows,

\[
(p_z - E)f(r) + \left( \frac{\partial}{\partial r} + \frac{m + \frac{1}{2}}{r} - \frac{1}{2} eB r \right) g(r) = 0,
\]

\[
\left( -\frac{\partial}{\partial r} + \frac{m - \frac{1}{2}}{r} - \frac{1}{2} eB r \right) f(r) + (-p_z - E) g(r) = 0,
\]

which are equivalent to

\[
\left\{ \frac{\partial^2}{\partial r^2} + \frac{1}{r} \frac{\partial}{\partial r} - \frac{(m - \frac{1}{2})^2}{r^2} - \left[p_z^2 - E^2 - eB \left(m + \frac{1}{2}\right)\right] - \frac{1}{4} e^2 B^2 r^2 \right\} f(r) = 0
\]
\[ g(r) = \frac{1}{p_z + E} \left( -\frac{\partial}{\partial r} + \frac{m - \frac{1}{2}}{r} - \frac{1}{2} eB r \right) f(r) \]  

We can define a dimensionless variable \( \rho = \frac{1}{2} eB r^2 \), then

\[ \frac{d}{dr} = eB r \frac{d}{d\rho}, \quad \frac{d^2}{dr^2} = eB \frac{d}{d\rho} + 2eB \rho \frac{d^2}{d\rho^2}. \]  

Now Eq. (A7) becomes

\[
\left\{ \frac{d^2}{d\rho^2} + \frac{d}{d\rho} - \left( m - \frac{1}{2} \right)^2 - \frac{1}{2eB} \left[ p_z^2 - E^2 - eB \left( m + \frac{1}{2} \right) \right] - \frac{1}{4} \rho \right\} f = 0
\]

Next, we choose

\[ f = e^{-\frac{\rho}{2}} \rho^{m+\frac{1}{2}} G(\rho), \]  

then Eqs. (A8, A10) become

\[ g = -\frac{\sqrt{2eB}}{E + p_z} e^{-\frac{\rho}{2}} \rho^m \frac{1}{2} G'(\rho), \]

\[ \rho G'' + \left[ \left( m + \frac{1}{2} \right) - \rho \right] G' - \frac{1}{2eB} \left( p_z^2 - E^2 \right) G = 0. \]

Define following two quantities,

\[ \gamma = m + \frac{1}{2}, \quad \alpha = \frac{1}{2eB} \left( p_z^2 - E^2 \right), \]

then Eq. (A13) becomes

\[ \rho G'' + (\gamma - \rho) G' - \alpha G = 0, \]

which is the confluent hypergeometric equation [33]. With the boundary conditions, \(|f(0)|, |f(\infty)| < \infty\), the solutions for \( G(\rho), f(\rho), g(\rho) \) can be chosen as:

1. When \( \gamma = 0, -1, -2, \ldots \), i.e. \( m = -1/2, -3/2, -5/2, \ldots \), the boundary condition \(|f(0)| < \infty\) requires that

\[ G(\rho) = \rho^{1-\gamma} F(\alpha - \gamma + 1, 2 - \gamma, \rho) = \rho^{\frac{1}{2} - m} F \left( \alpha - m + 1, \frac{3}{2} - m, \rho \right), \]

where \( F(\alpha, \gamma, \rho) \) is the confluent hypergeometric function as discussed in Appendix B. In addition, the boundary condition \(|f(\infty)| < \infty\) requires that

\[ \alpha - m + \frac{1}{2} = -n \quad (n \in \mathbb{N}), \quad E = \lambda \sqrt{p_z^2 + 2eB \left( n - m + \frac{1}{2} \right)} \quad (\lambda = \pm 1), \]

(1)
In addition, the boundary condition we can obtain the normalized eigenfunctions as follows:

\[ G(\rho) = \rho^{\frac{1}{2} - m} F\left(-n, \frac{3}{2} - m, \rho\right) \sim \rho^{\frac{1}{2} - m} L_n^{\frac{1}{2} - m}(\rho), \]

where \( L_n^{k}(\rho) \) is the general Laguerre polynomial as discussed in Appendix. Then one obtain \( f(\rho) \sim e^{-\frac{\rho}{2}} \rho^{\frac{1}{2} - \frac{m}{2}} L_n^{\frac{1}{2} - m}(\rho), \quad g(\rho) \sim -\frac{\sqrt{2eB}}{E + p_z} \left(n - m + \frac{1}{2}\right) e^{-\frac{\rho}{2}} \rho^{-\frac{1}{2} - \frac{m}{2}} L_n^{-\frac{1}{2} - m}(\rho). \) \hspace{1cm} (A19)

(2) When \( \gamma = 1, 2, 3, \ldots \), i.e. \( m = 1/2, 3/2, 5/2, \ldots \), the boundary condition \(|f(0)| < \infty \) requires that

\[ G(\rho) = F(\alpha, \gamma, \rho) = F\left(\alpha, \frac{1}{2} + m, \rho\right). \]

In addition, the boundary condition \(|f(\infty)| < \infty \) requires that

\[ \alpha = -n \quad (n \in \mathbb{N}), \quad E = \lambda \sqrt{p_z^2 + 2eBn} \quad (\lambda = \pm 1), \]

\[ G(\rho) = F\left(-n, \frac{1}{2} + m, \rho\right) \sim L_n^{m - \frac{1}{2}}(\rho). \]

Then one obtain

\[ f(\rho) \sim e^{-\frac{\rho}{2}} \rho^{\frac{m}{2} - \frac{1}{4}} L_n^{m - \frac{1}{2}}(\rho), \quad g(\rho) \sim \frac{\sqrt{2eB}}{E + p_z} e^{-\frac{\rho}{2}} \rho^{\frac{m}{2} - \frac{1}{4}} L_n^{m - \frac{1}{2}}(\rho). \] \hspace{1cm} (A23)

There is a special case we must point out here: When \( m > 0, n = 0 \), we must choose \( E = p_z \), in which case we have \( f(\rho) = e^{-\frac{\rho}{2}} \rho^{\frac{m}{2} - \frac{1}{4}}, \quad g(\rho) = 0 \). There is no physical solution for \( m > 0, n = 0, \quad E = -p_z \).

Making use of the orthonormal relation of the general Laguerre polynomials,

\[ \int_0^\infty dx e^{-x} x^\gamma L_m^\gamma(x) L_n^\gamma(x) = \frac{\Gamma(n + \gamma + 1)}{n!} \delta_{mn}, \]

we can obtain the normalized eigenfunctions as follows:

When \( m < 0 \),

\[ \psi_{\lambda nmpz} = \sqrt{\frac{n!}{(n - m + \frac{1}{2})!}} \left( \frac{\sqrt{eB(E + p_z)}}{2E} e^{-\frac{\rho}{2}} \rho^{\frac{m}{2} - \frac{1}{4}} L_n^{\frac{1}{2} - m} e^{i(m - \frac{1}{2})\phi} \right) \frac{e^{izp_z}}{2\pi}, \]

\[ E = \lambda \sqrt{p_z^2 + 2eB \left(n - m + \frac{1}{2}\right)}. \] \hspace{1cm} (A26)

When \( m > 0 \),

\[ \psi_{\lambda nmpz} = \sqrt{\frac{n!}{(n + m - \frac{1}{2})!}} \left( \frac{\sqrt{eB(E + p_z)}}{2E} e^{-\frac{\rho}{2}} \rho^{\frac{m}{2} + \frac{1}{4}} L_n^{\frac{1}{2} + m} e^{i(m + \frac{1}{2})\phi} \right) \frac{e^{izp_z}}{2\pi}, \] \hspace{1cm} (A27)
\[ E = \lambda \sqrt{p_x^2 + 2eBn}. \] (A28)

All normalized eigenfunctions are orthogonal with each other,

\[ \int dV \psi_{\lambda n'm'}^{\dagger} \psi_{\lambda m} = \delta_{\lambda \lambda'} \delta_{n'n} \delta_{m'm} \delta(p'_z - p_z). \] (A29)

### Appendix B: Confluent hypergeometric function and Laguerre polynomial

The confluent hypergeometric equation is \[33\]

\[ zy'' + (\gamma - z)y' - \alpha y = 0. \] (B1)

When \( \gamma \notin \mathbb{Z} \), there are two independent solutions as follows,

\[ y_1 = F(\alpha, \gamma, z), \]
\[ y_2 = z^{1-\gamma} F(\alpha - \gamma + 1, 2 - \gamma, z), \] (B2)

where \( F(\alpha, \gamma, z) \) is the confluent hypergeometric function defined as

\[ F(\alpha, \gamma, z) = \sum_{k=0}^{\infty} \frac{(\alpha)_k z^k}{(\gamma)_k k!} \equiv 1 + \frac{\alpha}{\gamma} z + \frac{\alpha(\alpha + 1)}{\gamma(\gamma + 1)} \frac{z^2}{2!} + \frac{\alpha(\alpha + 1)(\alpha + 2)}{\gamma(\gamma + 1)(\gamma + 2)} \frac{z^3}{3!} + \cdots. \] (B3)

The asymptotic behavior of \( F(\alpha, \gamma, z) \) as \( z \to \infty \) is the same as \( e^z \). When \( \alpha \) is a non-positive integer, then \( F(\alpha, \gamma, z) \) becomes a polynomial.

The general Laguerre polynomial \( L_\gamma^m(z) \) is defined from \( F(\alpha, \gamma, z) \) as follows \[34\],

\[ L_\gamma^m(z) = \frac{\Gamma(\gamma + n + 1)}{n!\Gamma(\gamma + 1)} F(-n, \gamma + 1, z) = \binom{\gamma + n}{n} F(-n, \gamma + 1, z), \] (B4)

where \( \gamma \in \mathbb{R} \) and \( n \in \mathbb{N} \). Laguerre polynomial \( L_\gamma^m(z) \) satisfies following differential equation

\[ zy'' + (\gamma + 1 - z)y' + ny = 0 \] (B5)

We can rewrite Eq. (B5) as a type of Sturm-Liouville equation,

\[ \frac{d}{dz} \left( z^{\gamma+1} e^{-z} \frac{dy}{dz} \right) + nz^\gamma e^{-z} y = 0, \] (B6)

which gives the orthogonality of \( L_\gamma^m(z) \),

\[ \int_0^\infty dz e^{-z} z^\gamma L_\gamma^m(z)L_\gamma^n(z) = \frac{\Gamma(n + \gamma + 1)}{n!} \delta_{mn}. \] (B7)
When $\gamma = 0$, then $L_n^\gamma(z)$ becomes the normal Laguerre polynomial $L_n(z)$,

$$L_n(z) = L_n^0(z) = F(-n, 1, z). \quad (B8)$$

[1] H. T. Elze, M. Gyulassy, and D. Vasak, Nucl. Phys. B 276, 706 (1986).
[2] D. Vasak, M. Gyulassy, and H. T. Elze, Annals Phys. 173, 462 (1987).
[3] J.-H. Gao, Z.-T. Liang, S. Pu, Q. Wang, and X.-N. Wang, Phys. Rev. Lett. 109, 232301 (2012), 1203.0725.
[4] S.-Z. Yang, J.-H. Gao, Z.-T. Liang, and Q. Wang, Phys. Rev. D 102, 116024 (2020), 2003.04517.
[5] X.-L. Sheng, R.-H. Fang, Q. Wang, and D. H. Rischke, Phys. Rev. D 99, 056004 (2019), 1812.01146.
[6] M. Buzzegoli, E. Grossi, and F. Becattini, JHEP 10, 091 (2017), [Erratum: JHEP 07, 119 (2018)], 1704.02808.
[7] M. Buzzegoli and F. Becattini, JHEP 12, 002 (2018), 1807.02071.
[8] A. Palermo, M. Buzzegoli, and F. Becattini, JHEP 10, 077 (2021), 2106.08340.
[9] V. E. Ambruş and E. Winstanley, Phys. Lett. B 734, 296 (2014), 1401.6388.
[10] D. Cangemi and G. V. Dunne, Annals Phys. 249, 582 (1996), hep-th/9601048.
[11] C. Zhang, R.-H. Fang, J.-H. Gao, and D.-F. Hou, Phys. Rev. D 102, 056004 (2020), 2005.08512.
[12] R.-H. Fang, R.-D. Dong, D.-F. Hou, and B.-D. Sun (2021), 2105.14786.
[13] A. Vilenkin, Phys. Lett. B 80, 150 (1978).
[14] A. Vilenkin, Phys. Rev. D 20, 1807 (1979).
[15] A. Vilenkin, Phys. Rev. D 22, 3080 (1980).
[16] M. E. Peskin and D. V. Schroeder, An introduction to quantum field theory, Westview Press, New York (1995).
[17] L. E. Parker and D. J. Toms, Quantum field theory in curved spacetime, Cambridge university Press, Cambridge (2009).
[18] H.-L. Chen, K. Fukushima, X.-G. Huang, and K. Mameda, Phys. Rev. D 93, 104052 (2016), 1512.08974.
[19] R.-D. Dong, R.-H. Fang, D.-F. Hou, and D. She, Chin. Phys. C 44, 074106 (2020), 2001.05801.
[20] D. E. Kharzeev, L. D. McLerran, and H. J. Warringa, Nucl. Phys. A 803, 227 (2008), 0711.0950.
[21] K. Fukushima, D. E. Kharzeev, and H. J. Warringa, Phys. Rev. D 78, 074033 (2008), 0808.3382.
[22] D. T. Son and P. Surowka, Phys. Rev. Lett. 103, 191601 (2009), 0906.5044.
[23] D. E. Kharzeev and D. T. Son, Phys. Rev. Lett. 106, 062301 (2011), 1010.0038.
[24] D. T. Son and N. Yamamoto, Phys. Rev. Lett. 109, 181602 (2012), 1203.2697.
[25] K. Landsteiner, E. Megias, and F. Pena-Benitez, Phys. Rev. Lett. 107, 021601 (2011), 1103.5006.
[26] S. Golkar and D. T. Son, JHEP 02, 169 (2015), 1207.5806.
[27] D.-F. Hou, H. Liu, and H.-c. Ren, Phys. Rev. D 86, 121703 (2012), 1210.0969.
[28] S. Lin and L. Yang, Phys. Rev. D 98, 114022 (2018), 1810.02979.
[29] J.-h. Gao, J.-Y. Pang, and Q. Wang, Phys. Rev. D 100, 016008 (2019), 1810.02028.
[30] A. Shitade, K. Mameda, and T. Hayata, Phys. Rev. B 102, 205201 (2020), 2008.13320.
[31] X.-l. Sheng, D. H. Rischke, D. Vasak, and Q. Wang, Eur. Phys. J. A 54, 21 (2018), 1707.01388.
[32] X.-L. Sheng, Ph.D. thesis, Frankfurt U. (2019), 1912.01169.
[33] J.-Y. Zeng, Quantum mechanics (Vol.1), Science Press, Beijing (2007).
[34] I. S. Gradshteyn and I. M. Ryzhik, Table of integrals, series, and products (Eighth edition), Academic Press, Oxford (2014).