Hawking radiation viewed as Landauer transport

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Viewing Hawking radiation as a 1D single quantum channel Landauer transport process, Nation et al calculated the energy flux and entropy flux from a Schwarzschild black hole without chemical potential. To generalize the method to the case with chemical potential, a rotating charged and non-charged BTZ black hole is investigated. Energy flux and entropy flux obtained are consistent with that from anomaly theory. The maximum energy flux and entropy flux are independent on the statistics of bosons or fermions.

Keywords: BTZ black hole; energy flux; gravitational anomaly; Landauer transport; fractional statistics

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I. INTRODUCTION

Hawking radiation, which is the most striking quantum effect arising from the quantum field in a curved spacetime background, attracts many efforts to investigate it. There are several kinds of derivations to obtain it. Hawking’s original derivation is calculating the Bogoliubov coefficients between the “in” and “out” states in a black hole background [1-3]. The subsequent Damour-Ruffini method is calculating the particles’ emitting rate by analytically extending the outgoing wave from outside of horizon to inside [4]. Parikh and Wilczek calculated WKB amplitude by considering pair creation of particles and antiparticles near the horizon and supposing that particles tunnel across the classical forbidden path where the barrier is created by the outgoing particles themselves [5-8]. Another approach to calculate Hawking radiation is due to anomaly through calculating the energy-momentum tensor in the black hole backgrounds [9-13]. The anomaly in field theory occurs if the symmetry of the action or the corresponding conservation law is valid in the classical theory but will violate in the quantized case. Anomalies can include conformal anomaly (or trace anomaly), anomaly in gauge symmetries, and gravitation anomaly. The pioneer work of Christensen and Fulling [9] told us that the strength of Hawking radiation flux is determined by the trace of energy-momentum tensor. Thinking of the two dimensional massless field, either Hawking effect or conformal anomaly can be deduced from the other. In Robinson and Wilczek’s work [10], Hawking radiation can be understood as compensating flux to cancel gravitation anomaly at the horizon.

In 80’s, Zurek has viewed Hawking radiation as a 3D black body radiation obeying Stefan-Boltzman law [31]. However, some recent works is indicating that a 4D black hole metric can be reduced to (1+1) dimensional spacetime and defined flat Rindler spacetime by virtue of the conformal symmetry near the horizon , thus Hawking radiation is inherently a (1+1) dimensional process. Recently, a new 1D single quantum channel transport model was used to explain Hawking radiation [14]. This model was first proposed to measure the electronic conductance of electrical transport in mesoscopic physics and was subsequently extended to thermal transport. It has been proved that the thermal conduction of 1D ballistic transport based on fractional statistics and the Landauer formulation is independent on the statistics nature and is governed by the universal quantum 1D quantum channel model is consistent with that from conformal symmetry arising near the horizon of (1+1)-dimensional Schwarzschild black hole. In Ref. [14], Hawking radiation energy and entropy flow of a Schwarzschild black hole is viewed as a 1D single quantum transport process. For a Schwarzschild black hole, the chemical potential $\mu_{BH} = 0$, the flux of bosons such as photons and gravitons is equal to the result of 1D quantum transport in the degenerate limit. For fermions such as neutrinos and electrons, in order to get the same maximum flux, the bi-direction current of particles and antiparticles must be taken into account. Meanwhile, the universal upper bound $S_{1D}^2 \leq \left( \frac{\pi^2}{3} \right) \dot{E}_{1D}$ holds all the time. The result obtained from 1D quantum channel model is consistent with that from conformal symmetry arising near the horizon of (1+1)-dimensional Schwarzschild black hole. In Refs. [19, 20], the method has been extended to some more complicated black holes with nonzero chemical potential and it is found that charge flux and gauge flux can also be viewed as a 1D Landauer transport process.

(2+1)-dimensional BTZ black hole as a solution of standard Einstein field equation $G_{\alpha\beta} + \Lambda g_{\alpha\beta} = \kappa T_{\alpha\beta}$ with negative cosmological constant $\Lambda = -l^2$ [21-24], has attracted great attention because it provides a simplified model for exploring black hole thermodynamics and quantum gravity. Hawking radiation from a BTZ black hole has been investigated in Refs. [25-27]. We will generalize Nation’s method to the case of nonzero chemical potential of the rotating charged BTZ black hole and a special case with $Q = 0$ in this paper.

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II. THE METRIC OF A BTZ BLACK HOLE

The line element of a charged rotating BTZ black hole can be written as \[ ds^2 = -N^2(r)dt^2 + N^{-2}(r)dr^2 + r^2[N\varphi(r)dt + d\varphi]^2, \] (1)
where the squared lapse \( N^2(r) \) and the angular shift \( N\varphi(r) \) are given as
\[ N^2(r) = -M + \frac{r^2}{l^2} + \frac{J^2}{4r^2} - \frac{\pi}{2}Q^2 \ln r \equiv f(r), \] (2)
\[ N\varphi(r) = -\frac{J^2}{2r^2}, \] (3)
with \(-\infty < t < \infty\), \(0 < r < \infty\), and \(0 \leq \varphi \leq 2\pi\). The black hole is characterized by the parameters ADM mass \( M \), angular momentum \( J \), and electronic charge \( Q \) carried by the black hole, which determine the asymptotic behavior of the solution. The metric is stationary and axially symmetric with Killing vectors \( \partial_t \) and \( \partial_\varphi \).

Horizons of the charged rotating BTZ black hole are roots of the lapse function. We only care the case while two distinct real roots exist, then the following inequality should be satisfied
\[ M > \frac{\pi Q^2 + \sqrt{\pi^2 Q^4 + \frac{16J^2}{l^2}}}{8} + \frac{2J^2}{f\left(\pi Q^2 + \sqrt{\pi^2 Q^4 + \frac{16J^2}{l^2}}\right)}. \] (4)

Hawking temperature is
\[ T_{BH} = \frac{1}{4\pi} f'(r)|_{r_-} = \frac{1}{2\pi r_-}\left(\frac{r_-^2}{l^2} - \frac{J^2}{2r_-^2} - \frac{\pi Q^2}{4}\right). \] (5)

Especially, \( Q = 0 \) is corresponding to a non-charged rotating BTZ black hole, and then the outer and inner event horizons \( r_\pm \) are given as
\[ r_\pm^2 = \frac{f}{2}\left(M \pm \sqrt{M^2 - \frac{J^2}{l^2}}\right). \] (6)
In terms of the inner and outer horizons, the black hole mass and angular momentum are
\[ M = \frac{r_+^2}{l^2} + \frac{J^2}{4r_+^2} = \frac{r_-^2 + r_+^2}{l^2}, \quad J = \frac{2r_+r_-}{l}, \] (7)
with the corresponding angular velocity
\[ \Omega = \frac{J}{2r_-^2}. \] (8)

Hawking temperature is
\[ T_{BH} = \frac{1}{4\pi} f'(r)|_{r_-} = \frac{1}{2\pi r_+}\left(\frac{r_-^2}{l^2} - \frac{J^2}{4r_+^2}\right). \] (9)

For the region near the horizon of a rotating BTZ black hole, the quantum field can be effectively described by an infinite collection of (1+1)-dimensional fields. The Kaluza-Klein reduction of the (2+1)-dimensional BTZ black hole yields
\[ ds^2 = -f(r)dt^2 + f^{-1}(r)dr^2, \] (10)
with a \( U(1) \) gauge field
\[ A_t = -\frac{J}{2r^2} - \frac{1}{2}Q^2 \ln r, \quad A_\varphi = -\frac{J}{2r^2}, \] (11)
and they are corresponding to the charged and uncharged circumstances respectively.
III. HAWKING RADIATION FLUX CALCULATED USING ANOMALIES

Hawking radiation from a rotating BTZ black hole can be obtained from gauge and gravitational anomalies [28]. For a reduced two-dimensional metric Eq. (10), the gravitational anomaly of the chiral scalar field is

\[ \nabla_{\mu} T_{\nu} = \frac{1}{96\pi} \frac{e^{\hat{\mu}}}{\sqrt{-g}} \partial_{\mu} \bar{\partial}_{\nu} \Gamma_{\bar{\nu}m}. \]  

(12)

which can also be rewritten as

\[ \nabla_{\mu} T_{\nu} \equiv A_{\nu} = \frac{1}{\sqrt{-g}} \partial_{\mu} N_{\nu}^{\mu}, \]  

(13)

where \( N_{\nu}^{\mu} \) is defined as

\[ N_{\nu}^{\mu} = \frac{1}{96\pi} e^{\hat{\mu}} \partial_{\nu} \partial_{\bar{\alpha}} \Gamma_{\bar{\alpha}m}. \]  

(14)

and the epsilon tensor reads

\[ \epsilon_{\mu\nu} = \left( \begin{array}{cc} 0 & 1 \\ -1 & 0 \end{array} \right). \]  

(15)

After calculating, non-zero components of \( N_{\nu}^{\mu} \) are

\[ N_{r}^{t} = \frac{1}{192\pi} \left( f'^2 + f'' f \right), \]  

(16)

\[ N_{t}^{r} = -\frac{1}{192\pi f^2} \left( f'^2 - f'' f \right). \]  

(17)

Now considering the gauge anomaly, neglecting classically irrelevant ingoing modes near the horizon, the effective two-dimensional theory becomes chiral near the horizon and the gauge symmetry becomes anomalous.

Outside far from the horizon, the current is conserved with

\[ \partial_r J_{(o)} = 0, \]  

(18)

while in the region near the horizon, since there is only outgoing field, the current satisfies the anomalous equation as

\[ \partial_r J_{(H)} = \frac{m^2}{4\pi} \partial_r A_{t}, \]  

(19)

where \( m \) is the \( U(1) \) charge.

By integrating Eq. (18) and Eq. (19), one can obtain the flux in each region as

\[ J_{(o)} = C_0, \quad J_{(H)} = C_H + \frac{m^2}{4\pi} [A_{t}(r) - A_t(r_+)]. \]  

(20)

where \( C_0 \) and \( C_H \) are integration constants. Under gauge transformations, variation of effective action is

\[ -\delta W = \int d^2x \sqrt{-g_2} d\nabla_{\mu} J^{\mu}, \]  

(21)

where \( \lambda \) is a gauge parameter and

\[ J^{\mu} = J_{(o)}^{\mu} \theta_+(r) + J_{(H)}^{\mu} H(r), \]  

(22)

where \( \theta_+(r) = \theta(r - r_+ - \epsilon), H(r) = 1 - \theta_+(r) \).

Using the anomaly equation, we get

\[ -\delta W = \int d^2x \lambda \left[ \delta(r - r_+ - \epsilon)(J_{(o)}^{\mu} - J_{(H)}^{\mu} + \frac{m^2}{4\pi} A_t) + \partial_r \left( \frac{m^2}{4\pi} A_t H \right) \right]. \]  

(23)
The total effective action should be gauge invariant and the last term should be canceled by quantum effects of the classically irrelevant ingoing modes. The quantum effect to cancel this term is induced by the ingoing modes near the horizon. The coefficient of the delta-function should vanish, so we have

\[ C_0 = C_H - \frac{m^2}{4\pi} A_t(r_+). \tag{24} \]

Since the covariant current is written as \( \tilde{J}^t = J^t + \frac{m^2}{4\pi} A_t H \), the condition \( \tilde{J}^t = 0 \) determines the value of the charge flux to be

\[ C_0 = -\frac{m^2}{4\pi} A_t(r_+). \tag{25} \]

Similarly the total flux of the energy-momentum tensor can be obtained as

\[ a_0 = \frac{m^2}{4\pi} A_t^2(r_+) + N_t'(r_+). \tag{26} \]

As for the charged rotating BTZ black hole with Eq. (2), the gauge potential \( A_t \) is given by Eq. (11) and \( N_t' \) is given by Eq. (16). Thus the gauge flux \( C_0 \) and total energy flux \( a_0 \) are written as

\[ C_0 = \frac{m^2}{4\pi} \left( \frac{J^t}{2} + \frac{1}{2} Q^2 \ln r_+ \right), \quad a_0 = \frac{m^2}{4\pi} \left( \frac{J^t}{2} + \frac{1}{2} Q^2 \ln r_+ \right)^2 + \frac{\pi}{12} T_{BH}^2. \tag{27} \]

Letting \( Q = 0 \) will give the case of uncharged BTZ black hole. We can find that gauge and gravitational anomalies in a BTZ black hole can be canceled by the total flux of Hawking radiation at Hawking temperature.

IV. LANDAUER TRANSPORT MODEL FOR HAWKING RADIATION FROM A BLACK HOLE

Now we will consider a single channel connecting two particle/heat reservoirs with (quasi-)particles obeying fractional statistics (generalized Bose and Fermi statistics) proposed by Haldane [29, 30]. The two reservoirs are characterized by the temperatures \( T_L \) and \( T_R \) with chemical potential \( \mu_L \) and \( \mu_R \), respectively. The subscripts \( L \) and \( R \) denote the left and right reservoirs respectively while we assume that \( T_L > T_R \) and the transport through 1D connection is adiabatic and ballistic (no scattering).

For particles obeying fractional statistics, the distribution function is given by

\[ f_g(x) = \frac{1}{\omega(x, g) + g}, \tag{28} \]

with \( \omega(x, g) \) given by the implicit function equation

\[ \omega^{\beta}(x, g)[1 + \omega(x, g)]^{-g} = e^x, \tag{29} \]

where \( x = \beta(E - \mu), \beta = \frac{1}{k_BT} \), \( g \) is the statistics parameter satisfying \( g \geq 0, g = 0 \) and \( g = 1 \) describe bosons and fermions respectively.

The net energy flux \( \dot{E} \) based on Landauer theory is \( \dot{E} = \dot{E}_L - \dot{E}_R \) with

\[ \dot{E}_{L(R)} = \frac{1}{\hbar} \int_{E_{L(R)}}^{E_{L(R)}^0} E f_g^{L(R)} dE, \tag{30} \]

where we have taken the canceling of the group velocity and density of state into account, and the particle transmission probability is supposed as 1.

Changing the integration variable from \( E \) to \( x = \beta(E - \mu) \), we have

\[ \dot{E}_{L(R)} = \frac{(k_B T_{L(R)})^2}{2\pi\hbar} \int_{\mu_{L(R)}}^{\mu_{L(R)}^0} dx \left( x + \frac{\mu_{L(R)}}{k_B T_{L(R)}} \right) f_g^{L(R)}(x). \tag{31} \]

Thinking of the fermion case, when the contribution of antiparticles is considered, the maximum energy flux of fermions is written as

\[ \dot{E}_{L(R)} = \frac{(k_B T_{L(R)})^2}{2\pi\hbar} \left[ \int_{\mu_{L(R)}}^{\mu_{L(R)}^0} dx \left( x + \frac{\mu_{L(R)}}{k_B T_{L(R)}} \right) \frac{1}{e^x + 1} + \int_{\mu_{L(R)}^0}^{\mu_{L(R)}^0} dy \left( y + \frac{\mu_{L(R)}}{k_B T_{L(R)}} \right) \frac{1}{e^{-y} + 1} \right]. \tag{32} \]
It can be rewritten as
\[
\dot{E}_{\text{L/R}} = \frac{(k_B T_{\text{L/R}})^2}{2\pi\hbar} \left[ \int_0^\infty dx \left( \frac{\mu_{\text{L/R}}}{k_B T_{\text{L/R}}(e^x + 1)} \right) + \int_0^\infty dx \left( -x + \frac{\mu_{\text{L/R}}}{k_B T_{\text{L/R}}(e^{-x} + 1)} \right) \right] \\
+ \int_0^\infty dy \left( y + \frac{\mu_{\text{L/R}}}{k_B T_{\text{L/R}}(e^y + 1)} \right) - \int_0^\infty dy \left( y + \frac{\mu_{\text{L/R}}}{k_B T_{\text{L/R}}(e^{-y} + 1)} \right)
\]
\[
= \frac{(k_B T_{\text{L/R}})^2}{2\pi\hbar} \left[ \int_0^\infty dx \frac{x}{e^x + 1} + \int_0^\infty dy \frac{y}{e^y + 1} + 2 \frac{\mu_{\text{L/R}}}{k_B T_{\text{L/R}}} \left( \int_0^\infty dx \frac{1}{e^x + 1} - \int_0^\infty dx \frac{1}{e^{-x} + 1} \right) + \frac{\mu_{\text{L/R}}^2}{2(k_B T_{\text{L/R}})^2} \right].
\]
so we have
\[
\dot{E} = \dot{E}_L - \dot{E}_R = \frac{\pi k_B^2}{12\hbar} (T_L^2 - T_R^2) + \frac{1}{4\pi\hbar} (\mu_L^2 - \mu_R^2),
\]
where we have considered that the upper limit of integral \(\frac{\mu_{\text{L/R}}}{k_B T_{\text{L/R}}(e^x + 1)}\) approaches to infinity.

For the charge flux, we have
\[
\dot{i} = \frac{k_B T_{\text{L/R}}e}{2\pi\hbar} \int_0^\infty dx \frac{1}{e^x + 1}.
\]
Considering the contribution of antiparticles, we get
\[
\dot{i} = \frac{k_B T_{\text{L/R}}e}{2\pi\hbar} \int_0^\infty dx \frac{1}{e^x + 1} + \frac{k_B T_{\text{L/R}}e}{2\pi\hbar} \int_0^\infty dy \frac{1}{e^y + 1}.
\]
Similarly, in the degenerate limit, the lower integration bound \(\frac{\mu_{\text{L/R}}}{k_B T_{\text{L/R}}(e^x + 1)}\) approaches to infinity, so we obtain
\[
\dot{l} = \frac{e}{2\pi\hbar}(\mu_L - \mu_R).
\]
As for the case of bosons, in the limit of degeneration, similar calculation as the fermion case can give the same conclusion as Eq. (33).

According to Ref. [14], the net entropy flux \(\dot{S}_{1D}\) is \(\dot{S}_{1D} = \dot{S}_L - \dot{S}_R\) with
\[
\dot{S}_{\text{L/R}} = -\frac{k_B T_{\text{L/R}}}{2\pi\hbar} \int_{\frac{\mu_{\text{L/R}}}{k_B T_{\text{L/R}}}}^\infty dx \left[ f_x \ln f_x + (1 - g f_x) \ln(1 - g f_x) - [1 + (1 - g) f_x] \ln [1 + (1 - g) f_x] \right].
\]
Changing integration variable \(x = \beta(E - \mu)\) to \(\omega\), the entropy flux can be simplified to
\[
\dot{S}_{\text{L/R}} = \frac{k_B^2 T_{\text{L/R}}}{2\pi\hbar} \int_{\omega_\beta}^\infty d\omega \frac{\ln(1 + \omega)}{\omega} - \frac{\ln \omega}{1 + \omega}.
\]
In the degenerate limit, the lower integration bound approaches to zero, therefore the statistics-dependence vanishes. The maximum entropy flux can be obtained
\[
\dot{S}_{1D} = \frac{\pi k_B^2}{6\hbar}(T_L - T_R).
\]
Till now, we have obtained the energy flux and charge flux for a 1D Landauer transport process. Taking Hawking radiation from a black hole as Landauer transport, where one reservoir is black hole with Hawking temperature \(T_L = T_{BH}\) and black hole’s electronic chemical potential \(\mu_L = \mu_{BH}\), the other reservoir is vacuum with \(T_R = \mu_R = 0\), we can give the total energy flux and charge flux as
\[
E = \dot{E}_L - \dot{E}_R = \frac{\pi k_B^2}{12\hbar} T_{BH}^2 + \frac{1}{4\pi\hbar} \mu_{BH}^2, \quad I = \frac{e}{2\pi\hbar} \mu_{BH},
\]
where \(\mu_{BH} = mA_l\). This is consistent with Eq. (27). The entropy flux is
\[
\dot{S}_{1D} = \frac{\pi k_B^2}{6\hbar} T_{BH}.
\]
In fact, when the 1D quantum transport system can be viewed as a near-equilibrium one, the electric flux ($\dot{I}$) and energy flux ($\dot{E}$) yield \[^{17}\]

$$\delta \dot{I} = \frac{\partial \dot{I}}{\partial \mu} \delta \mu + \frac{\partial \dot{I}}{\partial T} \delta T, \quad \delta \dot{E} = \frac{\partial \dot{E}}{\partial \mu} \delta \mu + \frac{\partial \dot{E}}{\partial T} \delta T. \quad (43)$$

with $\delta T = T_L - T_R$, $\delta \mu = \mu_L - \mu_R$.

For a system of fractional statistics under Eq. (28), taking the limit $\delta T \to 0$ and $\delta \mu \to 0$, the linear transport coefficients for arbitrary $g > 0$ can be given as \[^{17}\]

$$L_{11} = \frac{\partial \dot{I}}{\partial \mu} = M \frac{e}{2\pi \hbar} \int_0^\infty d\omega \frac{\omega}{(\omega + g)^2} = M \frac{e}{2\pi \hbar} \frac{1}{g}, \quad (44)$$

$$L_{12} = \frac{\partial \dot{I}}{\partial T} = M \frac{e}{2\pi \hbar} k_B \int_0^\infty d\omega \frac{\omega}{(\omega + g)^2} = 0, \quad (45)$$

$$L_{21} = \frac{\partial \dot{E}}{\partial \mu} = M \frac{1}{2\pi \hbar^2} \int_0^\infty d\omega \frac{\omega \log(\omega + g) + \mu \beta \omega}{(\omega + g)^2} = M \frac{\mu}{2\pi \hbar} \frac{1}{g}, \quad (46)$$

$$L_{22} = \frac{\partial \dot{E}}{\partial T} = M \frac{k_B^2 \pi^2}{2\pi \hbar^2} \int_0^\infty d\omega \frac{\omega^2 \log(\omega + g) + \mu \beta \omega}{(\omega + g)^2} = M \frac{k_B^2 \pi^2}{2\pi \hbar^2} \frac{\beta}{3}, \quad (47)$$

where $M$ is an integer related to the occupied modes number, we will neglect it by letting it to be 1.

So the energy flux and charge flux are

$$\delta \dot{E} = \frac{\mu}{2\pi \hbar g} \delta \mu + \frac{k_B^2 \pi^2}{2\pi \hbar^3} \delta T, \quad \delta \dot{I} = \frac{e}{2\pi \hbar g} \delta \mu. \quad (48)$$

It also means that

$$\dot{E}_{L,R} = \frac{\mu_{L,R}^2}{4\pi \hbar g} + \frac{\pi k_B^2 \hbar^2}{12\pi \hbar^3} \mu_{L,R}, \quad \dot{I} = \frac{e}{2\pi \hbar g} \mu. \quad (49)$$

V. CONCLUSION AND DISCUSSION

Following Nation's work, we have generalized it to the case with chemical potential. Viewing Hawking radiation as a 1D Landauer transport, we have calculated the energy flux and charge flux from a rotating charged and uncharged BTZ black hole. The total energy flux obtained from Landauer transport model, which is consistent with that from anomaly theory, contains not only thermal flux but also the contribution of flux caused by chemical potential. Based on the fractional statistics, the energy flux and entropy flux are independent on statistical behavior in the degenerate regime. From the above formulism Eqs. (48) and (49), we can easily find that the total energy flux of a 1D quantum transport system in the degenerate limit can be divided into two parts: the first term is due to the difference of chemical potential of the two reservoirs; the second term is purely thermal and entirely determined by the temperature. In addition, setting $\delta \mu = 0$, the total energy flux will eliminate to net thermal flux generated by $\delta T$. The 1D thermal conductance $\kappa^{\text{anom}} = \pi k_B^2 \hbar^2$ is exactly the coefficient $L_{22}$, which is independent on statistical behavior. As for the electronic flux, $L_{12} = 0$ means that it only depends on the chemical potential.

It is noticeable that Hawking radiation is viewed as a phenomenon near the horizon. As for infinity, the Hawking radiation would contain a gray-body factor, which caused by the effective potential outside the horizon.

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