Dissipation-Induced Luttinger Liquid Correlations in a One-Dimensional Fermi Gas

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We study a one-dimensional Fermi gas in the presence of dissipative coupling to environment through the Lindblad equation. The dissipation involves energy exchange with the environment and favours the relaxation of electrons to excitations. After switching on the dissipation, the system approaches a steady state, which is described by a generalized Gibbs ensemble. The fermionic single particle density matrix resembles deceivingly to that in a hermitian interaction quench. It decays inversely with the distance for short times due to the fermionic correlations in the initial state, which changes into a noninteger power law decay for late times, representing dissipation-induced Luttinger liquid behavior. However, the crossover between the two regions occurs due to dissipation-induced damping, and is unrelated to the propagation of excitations. The velocity of information spreading is set by the dissipative coupling, and differs significantly from the original sound velocity. The thermodynamic entropy grows as $-t \ln t$ initially, and saturates to an extensive value. Our results can be tested experimentally in one-dimensional Dirac systems.

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Introduction.—Thanks to the advent of sophisticated experimental technologies in cold atomic settings and in condensed matter, the creation and controlled manipulation of isolated quantum systems became possible [1–3]. In particular, one can follow the spatiotemporal dynamics of strongly interacting quantum gases [4,5] after some arbitrary time evolution protocol. The emerging universal behavior and scaling provides not only essential information on the (pre-)thermalization and relaxation, but is also relevant to simulate the early Universe after inflation, for which the experimental knobs are obviously limited. All this information becomes relevant for quantum computation and information processing [6].

However, no system is perfectly isolated from the environment, therefore considering open quantum systems, coupled to some external bath is necessary to understand realistic systems. In its simplest form, this is taken into account by the Lindblad equation [7–9]. This enterprise already gives way to engineer peculiar, dissipation-induced states of matter with no obvious analogues in closed quantum systems [10–16].

For closed quantum systems, Landau's Fermi liquid picture provides a good description of the normal state of many interacting systems in dimensions higher than one [17]. Therein, many properties of the original Fermi gas are inherited, though certain properties are renormalized. This picture breaks down in one dimension, and the ensuing interacting ground state differs markedly from that of the initial Fermi gas [18,19]. The original fermionic excitations are replaced by bosonic collective modes, consisting of many electron-hole pairs. Given the apparent vulnerability of a one-dimensional Fermi gas in closed quantum systems to Luttinger liquid or gap opening instabilities [18,19], their fate in an open quantum system is still an open question, i.e., when the fermionic degrees of freedom living in one dimension are coupled dissipatively to some environment. In particular, what is the structure of the ensuing steady state and what characterizes the nonunitary dynamical evolution towards the steady state?

This motivated us to investigate a one-dimensional Fermi gas in the presence of dissipative coupling to environment through the Lindblad equation. The dissipation involves energy exchange with the environment and favors the relaxation of electrons to excitations. We follow the full nonunitary dynamics of the system after switching on the dissipation at $t = 0$. Other systems were also investigated in similar context [20,21]. We find that the steady state is described *exactly* by a generalized Gibbs ensemble. The dissipation induces Luttinger liquid like correlation during the time evolution, but the velocity of information spreading is set by the dissipative coupling, and is unrelated to the

sound velocity. Our findings can be tested with current experimental technologies.

Dissipation in a one-dimensional Fermi gas.—We consider noninteracting one-dimensional spinless electrons, which, within the realm of a low energy theory, can propagate to left or right [18,19]. The low energy effective theory of the electrons in bosonized form gives rise to the Luttinger model with the Hamiltonian

$$H = \sum_{q>0} \omega_0(q)(b_q^+ b_q + b_{-q}^+ b_{-q}), \tag{1}$$

where b_q is the annihilation operator of the bosonic excitations which is bilinear of the original fermionic operators [18,19,22]. In Eq. (1), $\omega_0 = v|q|$ is the noninteracting spectrum with v the sound (or Fermi) velocity, describing low energy excitations around $\pm k_F$ with k_F the Fermi momentum.

In an open quantum system, coupling to environment induces nonunitary time evolution, which is described by the Lindblad equation as

$$\partial_t \rho = -i[H, \rho] + \sum_{q\neq 0} ([L_q, \rho L_q^+] + \text{H.c.}), \tag{2}$$

which determines the dynamics of the density matrix $\rho(t)$. This dissipative coupling to environment is taken into account by the jump operators L_q, which are best introduced visually in Fig. 1 for our current system. They only involve energy exchange with the environment and favor the relaxation of electrons to excitations. Similar jump operators were considered in Ref. [10]. The bosonized jump operators, visualized in Fig. 1 are

$$L_q = \sqrt{\gamma|q|}(\eta b_q + b_{-q}^+), \tag{3}$$

with $\eta > 0$ [23], and γ measures the strength of the coupling between the system and the environment and has velocity dimension. With this choice of jump operators, the ensuing problem becomes genuinely many body.

In general, the jump operators of the Lindblad equation describe the elementary processes occurring while the system interacts with its environment. The operators in Eq. (3) are chosen in such a way that they describe electron-hole excitation and relaxation while the total momentum of the system is shifted by momentum $-q$. Electron excitations increase the system energy with $\omega_0(q)$ while electron relaxation decreases it with the same amount. These processes are taken into account with different amplitudes and the parameter η enables us to describe either dissipation of energy to the environment or energy pumped into the system. Our choice of jump operators is also motivated by the possibility of studying dissipative effects analytically, focusing on features that do not depend qualitatively on the form of the coupling to the environment. Furthermore, the operator in Eq. (3) can be regarded a generalization of the electron density since L_q is proportional to the Fourier transformed electron density for $\eta = 1$. This limit was considered in Refs. [24,25].

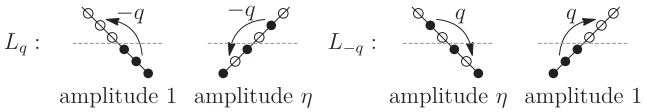

FIG. 1. Illustration of jump operators in the Lindblad equation, L_q and L_{-q} create electron-hole pairs on the right (around k_F) and left (around -k_F) moving branches with momentum $-q$ and q, respectively. The dashed line denotes the Fermi energy, and filled (empty) circles stand for electrons (holes). The jump operators are mixtures of excitation and relaxation of an electron with amplitude 1 and η, respectively.

To set the stage and to appreciate the role of η, we first calculate the time evolution of the average number of excitations, $\hat{n}_q = b_q^+ b_q$, as

$$\begin{aligned} n_q(t) &= \text{Tr}[\rho(t)\hat{n}_q] \\ &= \frac{1}{\eta^2 - 1} + \left(n_q(0) - \frac{1}{\eta^2 - 1}\right) e^{-2\gamma|q|t(\eta^2-1)}, \end{aligned} \tag{4}$$

where $n_q(0)$ is the occupation number in the initial state. For $\eta > 1$, i.e., when the boson annihilation has a larger amplitude compared to the boson creation, the boson number relaxes to $1/(\eta^2 - 1)$. This indicates that the system has a stable steady state. For $\eta \leq 1$, however, the occupation number explodes and the system is essentially boiled up to infinite temperatures.

By studying Eq. (2), it is remarkable that the Lindblad equation only couples q and $-q$ modes as long as the initial state does not couple additional modes. This allows us to write $\rho(t) = \prod_{q>0} \rho_q(t)$.

Time evolution and steady state of the Lindblad equation.—For one specific $q > 0$ mode, the solution of Eq. (2) is assumed in the form of

$$\begin{aligned} \rho_q(t) = r_q(t) e^{c_q(t) b_q b_{-q}} e^{-\ln[\nu_q(t)+1](b_q^+ b_q + b_{-q} b_{-q}^+)} \\ \times e^{c_q(t)^* b_q^+ b_{-q}^+}, \end{aligned} \tag{5}$$

where $\nu_q(t)$ and $r_q(t)$ are real functions of time and $c_q(t)$ is a complex-valued function. The trace of the density matrix is preserved when $\nu_q(t) > 0$ and

$$r_q(t) = \frac{\nu_q(t)^2 - |c_q(t)|^2}{\nu_q(t) + 1} > 0 \tag{6}$$

at any time instant. The latter equality shows that $r_q(t)$ is expressed with $\nu_q(t)$ and $c_q(t)$, therefore, the density matrix is completely characterized by these two functions. The average number of excitations is written as

$n_q(t) = \nu_q(t)/[\nu_q(t)^2 - |c_q(t)|^2]$ which has already been evaluated in Eq. (4).

By substituting Eq. (5) into the Lindblad equation Eq. (2), we obtain after some lengthy algebra [26]

$$\dot{\nu}_q = -2\gamma|q|[|c_q|^2 + \nu_q^2 + \nu_q(1-\eta^2) + \nu_q\eta(c_q + c_q^*)], \tag{7a}$$

$$\dot{c}_q = 2iv|q|c_q + 2\gamma|q|[c_q(\eta^2-1) - \eta(\nu_q^2 + c_q^2) - 2\nu_q c_q]. \tag{7b}$$

Despite their nonlinear nature, the differential equations can be solved and the stable steady states can be determined analytically [26]. For $\eta \leq 1$, the stable steady state is $\nu_{\mathrm{ex}} = 0$ and $c_{\mathrm{ex}} = 0$. For these values, however, no density matrix can be assigned since ν_{ex} is out of the domain $\nu > 0$. Nevertheless, the steady state can be interpreted physically as the boiled up system which is characterized by an infinite temperature. This is in accordance with the preliminary calculations of the occupation number. Namely, for $\eta \leq 1$, boson annihilation (electron relaxation) is not strong enough to damp the system, and the environment induces energy explosion.

For $\eta > 1$, when boson annihilation is expected to be strong enough to realize energy dissipation in the system, the stable fix point of the differential equations is

$$\nu_\infty = |A|^2\frac{\eta^2-1}{|A|^2-\eta^2} \qquad c_\infty = -\frac{\nu_\infty\eta}{|A|^2}A \tag{8}$$

with $A = 1 + [iv/\gamma(\eta^2-1)]$. Neither A nor the steady parameters ν_∞ and c_∞ depend on the wave number, therefore the stationary density matrix is the same in all $q > 0$ channels. This stationary density matrix is rewritten as

$$\rho_q(t\to\infty;\eta>1) = (1-e^{-\Omega_\infty})^2 e^{-\Omega_\infty(d_q^+ d_q + d_{-q}^+ d_{-q})}, \tag{9}$$

where $\Omega_\infty = |\mathrm{acosh}[(\nu_\infty^2 - |c_\infty|^2)/(2(\nu_\infty+1)) + 1]|$ is independent from the wave number $q > 0$. The operators d_q describe the eigenstates of the steady state and can be calculated via the Bogoliubov transformation [26]. The steady density matrix resembles a thermal state with $\omega_0(q)/T = \Omega_\infty$. The wave number independence of Ω_∞ implies that the temperature must depend on the momentum as $T(q) \sim |q|$. Therefore, the overall steady state density matrix describes exactly a generalized Gibbs ensemble (GGE) [30] [31]. Note that a similar density matrix describes only approximately the steady state of a Luttinger liquid after a hermitian interaction quench [33,34].

Single particle density matrix.—The density matrix enables us to calculate various physical quantities. We start with its single particle version, which is related to the Green's function in equilibrium. Since the fermion field decomposes to right and left moving parts as $\Psi(x) = e^{ik_F x}\Psi_R(x) + e^{-ik_F x}\Psi_L(x)$, it is enough [19,33] to concentrate on

$$G(x;t) = -i\mathrm{Tr}[\rho(t)\Psi_R^+(x)\Psi_R(0)] \tag{10}$$

where $\Psi_R(x)$ is the fermionic field operator of right-moving electrons with $\Psi_R(x) = (1/\sqrt{2\pi\alpha})\exp[i\sum_{q>0}\sqrt{2\pi/qL}(e^{iqx}b_q + e^{-iqx}b_q^+)]$, describing excitations around the right Fermi momentum k_F. By following standard steps [19,22], we obtain

$$\frac{G(x;t)}{G_0(x)} = \exp\left(-\sum_{q>0}\frac{4\pi}{Lq}n_q(t)[1-\cos(qx)]\right), \tag{11}$$

where L is the system size [35]. It is remarkable that all the time-dependence of the single particle density matrix occurs only through the average number of excitations. The function $G_0(x) = 1/(x+i\alpha)2\pi$ is the correlation function of noninteracting fermions at zero temperature. The length scale α is in the range of the lattice constant and is introduced as an exponential cutoff in momentum space, $\exp(-\alpha q)$.

The time-dependence of $n_q(t)$ is already obtained in Eq. (4). By starting initially from the ground state of the noninteracting Fermi gas, no excitations are present and $n_q(0) = 0$. In the thermodynamic limit, we obtain from Eq. (11)

$$\ln\frac{G(x;t)}{G_0(x)} = \frac{1}{1-\eta^2}\ln\left(\frac{1+(\frac{x}{\alpha})^2}{1+(\frac{x}{\alpha+2\gamma t(\eta^2-1)})^2}\right). \tag{12}$$

The most notable feature in Eq. (12) is that the time evolution is governed by the speed $\gamma(\eta^2-1)$, which can differ significantly from the original sound velocity v. For unitary time evolution (i.e., in the absence of dissipative coupling), any time dependence would be dictated by (a renormalized) v. For the Lindblad equation, however, the eigenvalues of the rhs of Eq. (2) have negative real part [7] (except for the steady state), whose magnitude is controlled by γ. After switching on the dissipation, any transient component of the density matrix dies out during the time evolution exactly due to the presence of γ. Therefore, the velocity of information spreading is set by the dissipative coupling, and differs from the original sound velocity.

At $t = 0$, the right-hand side of Eq. (12) vanishes and the correlation function is just equal to $G_0(x)$. After switching on the dissipative coupling at $t = 0$, the initial $G(x;t=0) \sim x^{-1}$ correlations are still retained for short times. Indeed, for $x \gg \gamma t(\eta^2-1)$, the Green's function gets dressed with a time dependent quasiparticle residue, $Z(t) = [1 + 2\gamma t(\eta^2-1)/\alpha]^{-2/(\eta^2-1)}$, which decays as a power law of time. This indicates that due to dissipation, the initial noninteracting state gets renormalized and heavy fermionic. On the other hand, for $x \ll \gamma t(\eta^2-1)$, the

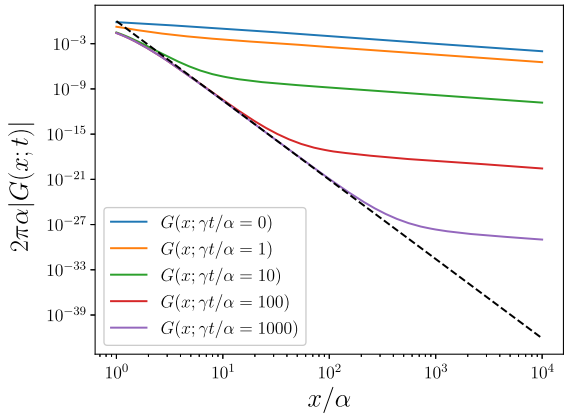

FIG. 2. The spatial dependence of the single particle density matrix, Eq. (12) for several times and $\eta = 1.3$. Initially at $\gamma t = 0$, it decays as x^{-1} (blue line). During the time evolution, the correlation function turns to the power-law decay as $x^{-(\eta^2+1)/(\eta^2-1)}$ for $x \ll \gamma t(\eta^2-1)$, and retains the x^{-1} decay for $x \gg \gamma t(\eta^2-1)$. The black dashed line denotes the steady state behavior.

fermionic nature of quasiparticles is lost and gives way to a noninteger Luttinger liquid like exponent, summarized as

$$G(x;t) \sim \begin{cases} \frac{Z(t)}{x} & \text{for } x \gg \gamma t(\eta^2-1) \\ (x/\alpha)^{-(\eta^2+1)/(\eta^2-1)} & \text{for } x \ll \gamma t(\eta^2-1) \end{cases}. \quad (13)$$

These features are highlighted in Fig. 2, indicating a smooth transition between the short and long distance decay.

The late time power-law exponent is $-(\eta^2+1)/(\eta^2-1)$. Similar noninteger exponents are familiar for Luttinger liquids, where the equilibrium exponent is well known [19] to be $-(K+K^{-1})/2$ at $T=0$ with K being the Luttinger parameter. In our setting, however, no electron interaction is present and the nontrivial exponent occurs solely due to the dissipation. It is important to note that the coupling to the environment does not lead to an effectively interacting *ground state* in the long time limit. On the contrary, we found that the steady density matrix rather represents a GGE which is a thermal state in general sense.

To elaborate on this a bit more, we consider the $2k_F$ oscillating part of the fermionic density correlation function [19,26]. For short times cf. Eq. (13), it retains its original fermionic spatial decay as $\sim\cos(2k_F x)/x^2 t^{\delta-2}$, while in the long time limit, it exhibits Luttinger liquid behavior in terms of noninteger spatial power law decay as $\sim\cos(2k_F x)/x^{\delta}$ with $\delta = 2[(\eta^2+1)/(\eta^2-1)] - 4\{[\eta(\eta^2-1)]/[(v/\gamma)^2+(\eta^2-1)^2]\} > 2$ [26]. Here, both characteristics of coupling to the environment, γ and η appear, while the single particle density matrix features only η. This indicates that the conventional Luttinger liquid relations between various exponents [19,36] do *not* hold for the dissipation-induced nontrivial steady state.

Entropy.—The relaxation toward the steady state is manifested also in the time evolution of the thermodynamic entropy, i.e., $S(t) = -\mathrm{Tr}[\rho(t)\ln\rho(t)]$. At $t=0$, the system is in a pure state with zero entropy. After switching on the coupling to the environment, the entropy varies with time as

$$S(t) = \sum_{q>0} \frac{2\Omega_q(t)}{e^{\Omega_q(t)}-1} - 2\ln\left(1-e^{-\Omega_q(t)}\right), \quad (14)$$

due to the nonunitary nature of the Lindblad equation. In Eq. (14), $\Omega_q(t)$ is the instantaneous eigenvalue of the exponent in the density matrix [26]. For weak dissipation, i.e., when $\gamma \ll v$ and $c_q(t) \approx 0$ is assumed in each momentum channel, the short time growth ($\gamma t \ll \alpha$) of the entropy is $S(t) \sim -L\gamma t\alpha^{-2}\ln(\gamma t/\alpha)$. The entropy satisfies a volume law and keeps growing for $\gamma t \sim \alpha$ and saturates to its steady value afterwards, which follows from substituting Ω_∞ into Eq. (14).

The single particle density matrix of our model exhibits similar time dependence to the evolution found after a quantum quench in the Luttinger model where the interaction was switched on suddenly [33,36], follows by unitary time evolution. Despite the similarities in the correlations, there are three essential differences: first of all, the velocity of information spreading in our case stems from the decay rate from Lindblad dynamics, while it originates from the propagation of quasiparticle excitations for the hermitian case and equals to the effective speed of light [37]. Second, the GGE is exact for the present dissipative system and is only approximate for the hermitian quantum quench [34,36]. Finally, dissipation leads to entropy production in our model while unitary time evolution does not change the entropy of the system.

Relation to experiments.—Experimentally, our setup can be realized by two coupled Luttinger liquids [38], interacting through chiral density-density interaction and without electron tunneling. One Luttinger liquid would represent the bath, which would directly induce the jump operators in Eq. (3) in the liquid, with $\gamma\eta^2$ and γ determined by the interaction between electron densities of the same and opposite chirality. In another setting, the jump operators can be implemented in a controlled fashion using a lattice realization of the Creutz ladder [39,40], which can also be realized experimentally [41]. When tuned to its critical point, it realizes one-dimensional Dirac fermions [42], and two legs of the ladder [42] host the right and left moving excitations, which are then also spatially separated. This allows for coupling the right and left moving densities to the environment independently to realize the dissipators depicted in Fig. 1. The unequal weights of the $\pm q$ processes in a given branch in Fig. 1 follow naturally, e.g., from the detailed balance [43].

Conclusion.—We have investigated the fate of a one-dimensional Fermi gas of electrons coupled to a dissipative

environment via the Lindblad equation. Using Abelian bosonization, the ensuing Lindblad dynamics is solved within the realm of the low energy effective theory. The steady state density matrix coincides with that of a generalized Gibbs ensemble. After switching on dissipation suddenly at $t = 0$, the single particle density matrix or Green's function exhibits similar spatiotemporal pattern than after a hermitian interaction quench [33]. This resemblance is, however, deceiving. Due to dissipation, correlations do not propagate with the effective sound velocity, but are damped by the dissipation, resulting in a significantly different velocity of information spreading. In addition, the characteristic features of Luttinger liquid correlation in terms of noninteger power law exponents for the spatial and temporal decay are revealed, but in this case these are induced by dissipation and not by electron-electron interaction. The thermodynamic entropy grows as $-t \ln t$ initially before saturating to its steady state value, and satisfies a volume law. These features can be observed in coupled Luttinger liquids or in one-dimensional Dirac systems.

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