Detecting Preformed-Pair Current through Nonequilibrium Noise in the BCS–BEC Crossover

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(Dated: February 9, 2022)

We theoretically propose a method to identify the tunneling current carrier in interacting fermions from nonequilibrium noise in the Bardeen-Cooper-Schrieffer to Bose–Einstein condensate crossover. The noise-to-current ratio, the Fano factor, can be a crucial probe for the current carrier. Bringing strongly-correlated fermions into contact with a dilute reservoir produces a tunneling current in between. The associated Fano factor increases from one to two as the interaction becomes stronger, reflecting the formation of the preformed Cooper pairs or bound molecules.

Introduction — Transport phenomena have contributed to the development of the fundamental physics in previous centuries. Various unconventional phenomena such as superfluidity and superconductivity were observed using transport measurements. However, clarifying the microscopic mechanism of the transport phenomena in strongly-correlated systems remains challenging because of their complexities such as strong interactions, lattice geometries, as well as multiple degrees of freedom.

Recently, an ultracold atomic system has been regarded as a quantum simulator for strongly-correlated many-body systems such as unconventional superconductors and nuclear systems, owing to its controllability of physical parameters (e.g., interparticle interactions and lattice structures) and its cleanliness, called atomtronics [1]. In particular, state-of-the-art experiments for tunneling current have been conducted in strongly interacting Fermi gases [2–5]. These experiments motivate us to study tunneling transport associated with the Josephson effect and Cooper-pair tunneling in the superfluid phase of the Bardeen-Cooper-Schrieffer (BCS) to Bose–Einstein-condensate (BEC) crossover [8–15].

One crucial problem is to determine how strong correlations affect transport quantities in the normal phase near the critical temperature (e.g., preformed pair and pseudogap [10]). Recently, several theoretical efforts have been paid to understand anomalous tunneling currents induced by pairing fluctuations in the normal phase [17–20], as observed in experiments [2–5]. It is reported that such anomalous pair currents can be induced by the nonlinear one-body tunneling processes [17], the one-body tunneling of a closed-channel molecule in the two-channel model [18], and the proximity effect associated with two-body interactions [21]. In this sense, it is worth exploring clear evidence for anomalous pair currents in a strongly interacting Fermi gas above the superfluid critical temperature.

For this purpose, measuring the Fano factor is promising, which is defined by a current and the associated nonequilibrium noise [22, 23]. The Fano factor in the large-biased setup reflects the effective charge per elementary transport process. The most fascinating example is the detection of fractional charges in fractional quantum Hall systems [24, 25]. The Fano factor has been used to determine the effective charge (or spin) in various physical systems such as superconductors [26, 27], Kondo quantum dots [28, 29], and magnetic junctions [30, 31].

In this study, we show that the Fano factor $F$ can be used as a probe for the current carrier in the BCS–BEC crossover. Figure 1 shows a schematic setup of the large-
biased system. Using the many-body $T$-matrix approach (TMA) \cite{34,35}, we numerically calculate the current and nonequilibrium noise within the Schwinger–Keldysh approach in the two-terminal tunneling junction under a large bias. Furthermore, we reveal how the Fano factor $F$ changes in a strongly-interacting regime, thereby reflecting the change of the dominant carrier. Our result can be tested by cold-atom experiments for which the noise measurement has been theoretically proposed \cite{36}.

Moreover, the Fano factor provides a direct evidence of pair-fluctuation effects associated with preformed pairs rather than other measurements such as spin susceptibility and photoemission spectra previously studied in this field \cite{40}. The current-noise measurement can also be used to identify the carriers of the BCS–BEC crossover in condensed-matter systems such as FeSe semimetal \cite{37,38}, lithium-intercalated layered nitrides \cite{41,42}, magic-angle twisted trilayer graphene \cite{43}, and organic superconductor \cite{44}. Moreover, the noise measurement has recently been conducted in a copper oxide heterostructure \cite{45,46} and disordered superconductor \cite{47}.

In the following, we take $\hbar = k_B = 1$ and consider a unit volume.

**Formalism**— We consider the Hamiltonian $\hat H = \hat H_L + \hat H_R + \hat H_{1T} + \hat H_{2T}$. The reservoir Hamiltonian $\hat H_{j=L,R}$ is given by

$$H_j = \sum_{p,\sigma} \xi_{p,j} c_{p,\sigma,j}^\dagger c_{p,\sigma,j} + g \sum_q P_{q,j}^\dagger P_{q,j},$$

(1)

where $\xi_{p,j} = p^2/(2m) - \mu_j$ denotes the kinetic energy measured from the chemical potential $\mu_j$ and $c_{p,\sigma,j}$ denotes the annihilation operator of a Fermi atom with momentum $p$ and the pseudospin $\sigma = \uparrow, \downarrow$. The second term in Eq. (1) denotes the attractive interaction with a contact-type coupling $g$, where $P_{q,j} = \sum_{p,q} \xi_{p,j} c_{p,q,j}^\dagger c_{p,q,j}$ is the pair-annihilation operator and $g$ is related to the scattering length $a$ as $g = \frac{\hbar^2}{m a} = \frac{1}{4} + \sum_p \frac{p^2}{8} \xi_{p,j}^2$. The one-body tunneling Hamiltonian, $H_{1T}$ is associated with the one-body potential barrier, where $t_{p,k}$ denotes its coupling strength. The two-body tunneling Hamiltonian reads

$$H_{2T} = \sum_{q,q'} \langle \omega_{q,\sigma,j}^2 \rangle P_{q,j}^\dagger P_{q',j}^\dagger P_{q,j} + \text{h.c.},$$

(3)

where $\omega_{q,q'}$ is the two-body coupling strength, induced by the local interaction term in Eq. (1) combined with the one-body tunneling of a closed-channel molecule in the two-channel model \cite{13} and the multiple one-body tunneling processes in the non-linear regime \cite{14,17,20,48}. Similar tunneling effects have also been examined in one-dimensional few-body systems \cite{49,50}. Here, we do not go into details on the origin of the one- and two-body tunneling, but rather investigate their possible consequence in observable quantities.

Using the Schwinger–Keldysh approach, we evaluate the expectation values of the current operator $\hat I = i [\hat N_L, \hat H]$ ($\hat N_j = \sum_{p,\sigma} c_{p,\sigma,j}^\dagger c_{p,\sigma,j}$ denotes the density operator in the $j$-reservoir) in the steady state at the lowest-order tunneling couplings by a sum of the one- and two-body contributions as $\hat I = \hat I_{qp} + \hat I_{pair}$, where each component reads \cite{21,51}.

$$I_{qp} = \frac{d\omega}{2\pi} \sum_{p,k,\sigma} \langle t_{k,p} \rangle^2 A_{k,L}(\omega) A_{p,R}(\omega) \times [f_L(\omega) - f_R(\omega)],$$

$$I_{pair} = 2 \frac{d\omega}{2\pi} \sum_{q,q'} \langle w_{q,q'} \rangle \times \langle B_{q,L}(\omega) B_{q',R}(\omega) \rangle \times [b_L(\omega) - b_R(\omega)].$$

(4)

In Eq. (4), $A_{k,j}(\omega)$ and $B_{q,j}(\omega)$ denote one- and two-particle spectral functions, respectively, $f_j(\omega)$ and $b_j(\omega)$ denotes the Fermi and Bose distribution functions, and $\mu_{s,j} = 2\mu_j$ denotes the bosonic-pair chemical potential in the $j$-reservoir. We define the current noise $S$ as

$$S = \frac{1}{T} \int_{-\infty}^{\infty} d\omega \left \langle \hat I(t) \hat I(0) \right \rangle.$$

(5)

Similar to the calculation above, we can evaluate the current noise \cite{51} as the sum of the two contributions: $S = S_{qp} + S_{pair}$, where

$$S_{qp} = \frac{d\omega}{2\pi} \sum_{p,k,\sigma} \langle t_{k,p} \rangle^2 A_{k,L}(\omega) A_{p,R}(\omega) \times [f_L(\omega) \{1 - f_R(\omega)\} + \{1 - f_L(\omega)\} f_R(\omega)],$$

$$S_{pair} = 4 \frac{d\omega}{2\pi} \sum_{q,q'} \langle w_{q,q'} \rangle \times \langle B_{q,L}(\omega) B_{q',R}(\omega) \rangle \times [b_L(\omega) \{1 - b_R(\omega)\} + b_R(\omega) \{1 + b_L(\omega)\}],$$

(6)

In the large bias limit ($\Delta \mu \equiv \mu_L - \mu_R \to \infty$), we can prove $S_{qp}/I_{qp} = 1$ and $S_{pair}/I_{pair} = 2$ without any further approximations \cite{51}. This motivates us to consider the Fano factor,

$$F = \frac{S}{I} = \frac{S_{qp} + S_{pair}}{I_{qp} + I_{pair}}.$$
given densities $N_1$ in the BCS–BEC crossover regime. The single-particle propagator is given by

$$G_{k,j}(i\omega_n) = \frac{1}{C_{k,j}^0(i\omega_n) - \Sigma_{k,j}(i\omega_n)},$$

(8)

$$\Sigma_{k,j}(i\omega_n) = T_j \sum_{q,\ell} \Gamma_{q,j}(i\nu_\ell)G_{q-k,j}(i\nu_\ell - i\omega_n),$$

(9)

where $C_{k,j}^0(i\omega_n) = (i\omega_n - \xi_{k,j})^{-1}$ denotes the bare propagator and $\Sigma_{k,j}(i\omega_n)$ denotes the TMA self-energy. Following a standard TMA procedure [52], the $T$-matrix $\Gamma_{q,j}(i\nu_\ell)$ is formulated by incorporating the particle–particle multiple scattering as

$$\Gamma_{q,j}(i\nu_\ell) = g [1 - g\Pi_{q,j}(i\nu_\ell)]^{-1},$$

(10)

using the bare two-body propagator given as

$$\Pi_{q,j}(i\nu_\ell) = -T_i \sum_{p,n} C_{p+q/2,j}^0(i\omega_n + i\nu_\ell)G_{-p+q/2,j}^0(-i\omega_n).$$

(11)

The fermion (boson) Matsubara frequency is denoted by $\omega_n$. Furthermore, we define the dressed two-body propagator [53] as

$$G_{q,j}(i\nu_\ell) = \Pi_{q,j}(i\nu_\ell) [1 + \Pi_{q,j}(i\nu_\ell)\Gamma_{q,j}(i\nu_\ell)].$$

(12)

The spectral functions can be obtained from the analytic continuation as $A_{k,j}(\omega) = -2\text{Im} G_{k,j}(i\omega_n - \omega - \mu_{1,j} + i\eta)$ and $B_{q,j}(\omega) = -2\text{Im} G_{q,j}(i\nu_\ell - \omega - \mu_{2,j} + i\eta)$ with an infinitesimal small number $\eta$.

In this study, we consider the large bias regime (see Fig. 1) characterized by $\mu_L - \mu_R \to \infty$ [51, 54] and the momentum-conserved tunneling processes as $\delta_{p,k} = T_{1,p,k}$ and $\omega_{q,q'} = T_{2,q,q'}$, for simplicity. We employ $\eta = 10^{-2}E_{F,L}$ in the numerical calculation to avoid the divergent behavior of the current associated with the momentum-conserved tunneling in the weak- and strong-coupling limits, where $E_{F,L} = (3\pi^2N_L)^{1/3}/(2m)$ denotes the Fermi energy of the $L$ reservoir with the number density $N_L$. However, our result can be qualitatively unchanged by this treatment because the distribution functions play a key role in determining $F$ rather than the detailed structures of tunneling junctions. Moreover, $T_2$ must be normalized to suppress the ultraviolet divergence in $B_{q,j}(\omega)$. For this purpose, we introduce the renormalized two-body tunneling $T_{2,\text{ren.}} = \frac{k_{F,L}}{\Lambda^2} T_2$, where $k_{F,L} = \sqrt{2mE_{F,L}}$ denotes the Fermi momentum. Such a divergence can also be avoided by introducing the form factor for the relative momentum $p$ in $P_{q,j}$. In this work, we take $\Lambda = 100k_{F,L}$ [53] in the practical calculation.

Results—Fig. 2 shows the Fano factor $F$, as a function of the dimensionless interaction parameter $(k_{F,L}a)^{-1}$ in the entire BCS-BEC crossover regime above the superfluid critical temperature $T_c$. We considered $T_{2,\text{ren.}}/T_1 = 1$, and the reservoir $R$ was regarded as almost vacuum ($\mu_L - \mu_R \to \infty$). One can clearly see that $F$ evolves from 1 to 2 with increasing the interaction strength in Fig. 2, indicating that the current carrier gradually changes from quasiparticles ($F = 1$) to preformed Cooper pair or bound molecules ($F = 2$). Such a behavior is universal in the sense that these asymptotic values do not depend on any details on the model parameters and structures of tunneling junctions. More explicitly, at the large bias limit, one can obtain

$$F(\Delta \mu \to \infty) \rightarrow \frac{I_{\text{qp}} + 2I_{\text{pair}}}{I_{\text{qp}} + I_{\text{pair}}},$$

(13)

where $I_{\text{qp}}$ and $I_{\text{pair}}$ denote the contributions of the quasiparticle and pair tunnelings, respectively. The Fano factor $F$ approaches 1 and 2 in the quasiparticle-dominant ($I_{\text{qp}} \gg I_{\text{pair}}$) and pair-dominant regimes ($I_{\text{pair}} \gg I_{\text{qp}}$), respectively. Although the interaction dependence of the Fano factor $F$ is deeply related to properties of the tunneling junctions and spectral functions of the carriers, one can find from Eq. 13 that $F \to 1$ ($F \to 2$) in the limit of $a^{-1} \to -\infty$ ($a^{-1} \to \infty$) regardless of the detailed properties of the system. Moreover, $F = 2$ can be realized even above $T_c$ because of strong interactions leading to the formation of preformed Cooper pairs in the BCS–BEC crossover. With increasing the temperature, $F$ tends to be suppressed because thermal effects assist the dissociation of pairs. Nevertheless, even at finite temperature, $F$ approaches 2 with increasing the interaction because bound molecules are dominant in the deep BEC regime where $T_L \lesssim E_b$ [$E_b = 1/(ma^2)$ is the two-body binding energy].

To see the detailed behavior of the Fano factor $F$, we
plot $I_{\text{qp}}$ and $I_{\text{pair}}$ throughout the BCS-BEC crossover at different temperatures in Fig. 3. From the inset of Fig. 3, the quasiparticle current $I_{\text{qp}}$ is exponentially suppressed with increasing the attractive interaction $(k_{F,L},a)^{-1} > 0$. This suppression is induced by the pseudogap effect, i.e., the reduction of $A_{k_{F,L}}(\omega)$ near $|k| = k_{F,L}$ and $\omega = E_{F,L} \approx \mu_L$ by the particle-hole coupling. Finally, $I_{\text{qp}}$ approaches zero in the BEC limit $(k_{F,L},a)^{-1} \rightarrow \infty$ because of the formation of molecules with large binding energies. These results are qualitatively consistent with previous work \cite{17, 20}. On the other hand, $I_{\text{pair}}$ drastically increases with increasing the interaction strength $(k_{F,L},a)^{-1}$ as shown in Fig. 3. At the BCS side $(k_{F,L},a)^{-1} < 0$ where the attraction is not strong to form a two-body bound state in vacuum, the contribution of $I_{\text{pair}}$ can be regarded as the tunneling of the preformed Cooper pairs into the two-body continuum in the reservoir R. In the strong-coupling BEC regime $((k_{F,L},a)^{-1} > 1$ and $T_L/E_b \lesssim 1)$, $I_{\text{pair}}$ describes the tunneling transport of bound molecules across two reservoirs, because the two-body bound state exists in the reservoir R with the same coupling $g$. Such a tunneling current associated with weakly-interacting molecular bosons becomes large due to their long lifetime and the Bose enhancement of low-energy distributions.

One can also see a dip-hump structure of $I_{\text{pair}}$ in the intermediate regime. Here, $\mu_L$ is close to zero and changes its sign, indicating that the dominant contribution changes from the preformed-pair transfer to the molecule-to-molecule transport across the junction. From the unitary limit $((k_{F,L},a)^{-1} = 0)$, the preformed-pair transfer increases due to the overlap with the bound-state spectra in $B_{q,R}(\omega)$ and eventually decreases because of the decrease in $\mu_L$. With increasing the interaction further, the inter-reservoir molecule-to-molecule transition emerges where the bound-state spectra in two reservoirs get close to each other in the energy axis $\omega$. Although these structures reflect the physical properties of the system, they also depend on the detailed setup of the tunneling junctions (e.g., the ratio between the tunneling couplings $T_{2\text{ren.}}/T_1$ \cite{51}.

Figure 4 shows the temperature dependence of the Fano factor $F$ in the unitary limit $((k_{F,L},a)^{-1} = 0)$. Because $B_{q,R}(\omega)$ does not involve a bound molecule pole, the transfer of the preformed Cooper pairs in the reservoir L to the scattering two-body continuum in the reservoir R can be anticipated in the unitary limit. One can see the enhancement of the Fano factor $F$ at the low-temperature regime. In particular, the curvature of the Fano factor $F$ is modified at $T_L/T_c \simeq 2.8$, where the sign of $\mu_L$ changes from negative to positive one as the temperature decreases (see the inset of Fig. 4). At a positive $\mu_L$, the pole of the preformed Cooper pairs gradually appears in $B_{q,L}(\omega)$. Thus, the behavior of the Fano factor $F$ can be regarded as a signature of the emergence of the preformed Cooper pairs. Because the preformed Cooper pairs play an important role in the pseudogap physics of ultracold Fermi gases \cite{16}, the Fano factor contributes to the further understanding of pairing pseudogaps in the BCS–BEC crossover regime. Incidentally, because TMA does not capture the self-energy shift in $\Pi_{q,L}(\omega)$, the curvature change of the Fano factor $F$ may differ from the temperature where $\mu_L = 0$ in actual experiments and in more sophisticated theoretical approaches \cite{34, 35}.

**Summary**—In this study, we showed that the Fano factor (i.e., the noise-to-current ratio $F = S/I$) can be a useful probe for current carriers in the BCS–BEC crossover at large-biased tunneling junctions. Using the many-body TMA, we demonstrated that the Fano factor...
where the dominant current carrier changes from the quasiparticle ($F = 1$) to the pair ($F = 2$) along the BCS-BEC crossover. Our prediction can be tested by experiments and uncover nonequilibrium strong-coupling physics via transport measurements. Furthermore, our result indicates that the noise measurement is useful for the study of the BCS-BEC crossover and pair-fluctuation effects in unconventional superconductors.

This work is supported in part by Grants-in-Aid for Scientific Research from JSPS (Grants Nos. JP18H05406, JP20K03831). D.O. is funded by the President’s PhD Scholarships at Imperial College London. MM is partially supported by the Priority Program of the Chinese Academy of Sciences, Grant No. XDB28000000.

Schwinger-Keldysh approach for current and noise

We start from the current operator given by

$$I = \hat{I}_{\text{qp}} + \hat{I}_{\text{pair}},$$

$$\hat{I}_{\text{qp}} = i \sum_{p,k,\sigma} t_{k,p} \left[ c_{k,\sigma,L}^\dagger c_{p,\sigma,R} - c_{k,\sigma,R}^\dagger c_{p,\sigma,L} \right],$$

$$\hat{I}_{\text{pair}} = 2i \sum_{q,q'} w_{q,q'} \left[ P_{q,L}^1 P_{q',R}^1 - P_{q',R}^1 P_{q,L}^1 \right],$$

where $\hat{I}_{\text{qp}}$ and $\hat{I}_{\text{pair}}$ are operators for quasiparticle and pair currents, respectively. Truncating the higher-order contributions with respect to the tunneling Hamiltonians [i.e., $O(H_3^3)$, $O(H_2^3)$], we can evaluate their expectation values, $I_{\text{qp}}(t_1, t_2) = \langle \Psi(t_1) | \hat{I}_{\text{qp}} | \Psi(t_2) \rangle$ and $I_{\text{pair}}(t_1, t_2) = \langle \Psi(t_1) | \hat{I}_{\text{pair}} | \Psi(t_2) \rangle$, for the different times $t_1$ and $t_2$, where $| \Psi(t) \rangle$ is the state-vector of the steady state. First, the quasiparticle contribution reads

$$I_{\text{qp}}(t_1, t_2) = -2 \int_C dt' \sum_{p,k,\sigma} |t_{k,p}|^2 \mathrm{Re} \left[ \langle T_C c_{k,\sigma,R} c_{p,\sigma,R}^\dagger (t_2) c_{k,\sigma,R} (t') \rangle \langle T_C c_{p,\sigma,R}^\dagger (t') c_{p,\sigma,L} (t_1) \rangle \right],$$

(17)

where $C$ denotes the Keldysh contour. Note that while the right hand side of Eq. (17) depends only on $t_1 - t_2$ in considering the steady state. Using the Green’s functions, we rewrite $I_{\text{qp}}(t_1, t_2)$ as

$$I_{\text{qp}}(t_1, t_2) = 2 \int_{-\infty}^\infty dt' \sum_{p,k,\sigma} |t_{k,p}|^2 \mathrm{Re} \left[ G_{p,R}^{\text{ret}}(t_2 - t') \langle G_{k,L}^\lt (t' - t_1) + G_{p,R}^\lt (t_2 - t') G_{k,L}^{\text{adv.}} (t' - t_1) \rangle \right],$$

(18)

where $G_{p,R}^{\text{ret.}(\text{adv.})}$ is the retarded (advanced) Green’s function of a fermion in thermal equilibrium. The lesser component $G^\lt$ contains the information of the thermal distribution in each reservoir. Here, we take $t_1 = t_2 \equiv t$ and the Fourier transformation

$$I_{\text{qp}} = 2 \int \frac{d\omega}{2\pi} \sum_{p,k,\sigma} |t_{k,p}|^2 \mathrm{Re} \left[ G_{p,R}^{\text{ret}}(\omega) G_{k,L}^\lt (\omega) + G_{p,R}^\lt (\omega) G_{k,L}^{\text{adv.}} (\omega) \right].$$

(19)

Moreover, we use

$$G_{k,j}^\lt (\omega) = -2 if_j(\omega) \Im G_{k,j}^{\text{ret}} (\omega) \equiv if_j(\omega) A_{k,j} (\omega),$$

(20)

where

$$f_j(\omega) = \frac{1}{\exp(\frac{\omega - \mu_j}{T_j}) + 1}$$

(21)

is the Fermi-Dirac distribution function. We use Matsubara Green’s functions in each reservoir reaching thermal equilibrium as a grand-canonical ensemble with $-\mu_j \bar{N}_j$ and obtain the retarded(advanced) Green’s function by the analytic continuation with $\mu_j$ as $i\omega_n \rightarrow \omega + i\eta - \mu_j$ in each reservoir. Then, we obtain

$$I_{\text{qp}} = \int \frac{d\omega}{2\pi} \sum_{p,k,\sigma} |t_{k,p}|^2 A_{p,L}(\omega) A_{k,R}(\omega) [f_L(\omega) - f_R(\omega)].$$

(22)
Similarly, we obtain the pair current contribution as

\[ I_{\text{pair}} = 2 \sum_{q,q'} \int \frac{d\omega}{2\pi} \omega |q_{q',q}|^2 B_{q,L}(\omega) B_{q'R}(\omega) [b_L(\omega) - b_R(\omega)], \]  

(23)

where we used the relation for the two-particle Green’s function given \( G^< \) by

\[ G_{q,j}^<(\omega) = 2ib_j(\omega) \text{Im} G_{q,j}^{\text{ret.}}(\omega) = -ib_j(\omega) B_{q,j}(\omega), \]  

(24)

and the Bose-Einstein distribution function

\[ b_j(\omega) = \frac{1}{\exp \left( \frac{\omega - \mu_j}{T} \right) - 1}, \]  

(25)

with a bosonic (pair) chemical potential \( \mu_{q,j} = 2\mu_j \). \( G^{<(>)} \) and \( G^{\text{ret. (adv.)}} \) are the lesser (greater) and retarded (advanced) components of two-particle Green’s functions, respectively. One can find that \( I = I_{\text{qp}} + I_{\text{pair}} \) obtained from Eqs. (17) and (23) is equivalent to Eq. (11).

Next, we consider the current noise

\[ S = \frac{1}{2} \int_{-\infty}^{\infty} dt \left( \langle \dot{I}(t) \dot{I}(0) \rangle + \langle \dot{I}(0) \dot{I}(t) \rangle \right). \]  

(26)

At lowest order of tunneling couplings, we obtain

\[ \langle \dot{I}(t) \dot{I}(0) \rangle = \sum_{p,k,\sigma} |t_{p,k,\sigma}|^2 \left[ G_{k,L}^<(t) G_{p,R}^>(-t) + G_{p,R}^<(t) G_{k,L}^>(-t) \right] \]
\[ - 4 \sum_{q,q'} |w_{q,q'}|^2 \left[ G_{q,L}^<(t) G_{q',R}^>(-t) + G_{q',R}^<(t) G_{q,L}^>(-t) \right], \]  

(27)

\[ \langle \dot{I}(0) \dot{I}(t) \rangle = \sum_{p,k,\sigma} |t_{p,k,\sigma}|^2 \left[ G_{k,L}^<(t) G_{p,R}^>(t) + G_{p,R}^<(t) G_{k,L}^>(t) \right] \]
\[ - 4 \sum_{q,q'} |w_{q,q'}|^2 \left[ G_{q,L}^<(t) G_{q',R}^>(t) + G_{q',R}^<(t) G_{q,L}^>(t) \right]. \]  

(28)

Collecting them and taking the Fourier transformation, we obtain

\[ S = S_{\text{qp}} + S_{\text{pair}}, \]  

(29)

\[ S_{\text{qp}} = \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \sum_{k,p,\sigma} |t_{k,p,\sigma}|^2 \left[ G_{k,L}^<(\omega) G_{p,R}^>(\omega) + G_{k,L}^>(\omega) G_{p,R}^<(\omega) \right], \]
\[ S_{\text{pair}} = -4 \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \sum_{q,q'} |w_{q,q'}|^2 \left[ G_{q,L}^<(\omega) G_{q',R}^>(\omega) + G_{q',R}^<(\omega) G_{q,L}^>(\omega) \right]. \]  

(30)

Using the relations associated with greater Green’s functions

\[ G_{p,j}^>(\omega) = -i A_{p,j}(\omega) [1 - f_j(\omega)], \quad G_{q,j}^>(\omega) = -i B_{q,j}(\omega) [1 + b_j(\omega)], \]  

(31)

and the lesser ones given by Eqs. (20) and (24), we obtain

\[ S_{\text{qp}} = \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \sum_{k,p,\sigma} |t_{k,p,\sigma}|^2 A_{k,L}(\omega) A_{p,R}(\omega) \left[ f_L(\omega) \{ 1 - f_R(\omega) \} + \{ 1 - f_L(\omega) \} f_R(\omega) \right], \]
\[ S_{\text{pair}} = 4 \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \sum_{q,q'} |w_{q,q'}|^2 B_{q,L}(\omega) B_{q',R}(\omega) \left[ b_L(\omega) \{ 1 + b_R(\omega) \} + b_R(\omega) \{ 1 + b_L(\omega) \} \right], \]  

(32)

which is equivalent to Eq. (12).
For a small bias limit at equal temperatures $T_L = T_R \equiv T$ where $\Delta \mu \rightarrow 0$ and $f_R(\omega) \rightarrow f_L(\omega) \equiv f(\omega)$ with $\mu_R \rightarrow \mu_L \equiv \mu$, we obtain

$$f_L(\omega) - f_R(\omega) = -\frac{\partial f(\omega)}{\partial \omega} \Delta \mu + O((\Delta \mu)^2), \quad (33)$$

$$b_L(\omega) - b_R(\omega) = -2 \frac{\partial b(\omega)}{\partial \omega} \Delta \mu + O((\Delta \mu)^2). \quad (34)$$

Using

$$f(\omega)\{1 - f(\omega)\} = -T\frac{\partial f(\omega)}{\partial \omega}, \quad b(\omega)\{1 + b(\omega)\} = -T\frac{\partial b(\omega)}{\partial \omega}, \quad (35)$$

we recover the Onsager’s relation

$$S(\Delta \mu \rightarrow 0) = 2T \frac{I}{\Delta \mu}. \quad (36)$$

Moreover, the current and the noise can be rewritten as

$$I_{qp} = \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \sum_{p,k,\sigma} |t_{k,p}|^2 A_{k,L}(\omega) A_{p,R}(\omega) \left[ -\frac{1}{2} \sinh \left( \frac{\beta_L(\omega - \mu_L) - \beta_R(\omega - \mu_R)}{2} \right) \right], \quad (37)$$

$$I_{\text{pair}} = 2 \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \sum_{q,q'} |w_{q,q'}|^2 B_{q,L}(\omega) B_{q',R}(\omega) \left[ -\frac{1}{2} \sinh \left( \frac{\beta_L(\omega - \mu_{L,R}) - \beta_R(\omega - \mu_{L,R})}{2} \right) \right], \quad (38)$$

$$S_{qp} = \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \sum_{k,p,\sigma} |t_{k,p}|^2 A_{k,L}(\omega) A_{p,R}(\omega) \left[ -\frac{1}{2} \cosh \left( \frac{\beta_L(\omega - \mu_L) - \beta_R(\omega - \mu_R)}{2} \right) \right], \quad (39)$$

$$S_{\text{pair}} = 4 \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \sum_{q,q'} |w_{q,q'}|^2 B_{q,L}(\omega) B_{q',R}(\omega) \left[ -\frac{1}{2} \cosh \left( \frac{\beta_L(\omega - \mu_{L,R}) - \beta_R(\omega - \mu_{L,R})}{2} \right) \right]. \quad (40)$$

In particular, considering the large-biased limit where

$$\tanh \left( \frac{\beta_L(\omega - \mu_L) - \beta_R(\omega - \mu_R)}{2} \right) \approx \tanh \left( \frac{\beta_L(\omega - \mu_{L,R}) - \beta_R(\omega - \mu_{L,R})}{2} \right) \approx 1, \quad (41)$$

is satisfied, we obtain

$$S_{qp}(\Delta \mu \rightarrow \infty) \rightarrow I_{qp}, \quad S_{\text{pair}}(\Delta \mu \rightarrow \infty) \rightarrow 2I_{\text{pair}}, \quad (42)$$

where we have denoted $I \equiv I_{qp} + I_{\text{pair}}$. The result of Eq. [42] motivates us to consider the Fano factor

$$F = \frac{S}{I} = \frac{S_{qp} + S_{\text{pair}}}{I_{qp} + I_{\text{pair}}}. \quad (43)$$

Then, one can see that the Fano factor $F$ in a large-biased junction changes from 1 to 2 reflecting the ratio between $I_{qp}$ and $I_{\text{pair}}$.

**RETARDED PROPAGATORS IN THE DILUTE RESERVOIR**

For the single-particle Green’s function in the reservoir $R$ at dilute limit, we employ the non-interacting one given by

$$G_{p,R}^{\text{ret}}(\omega) = \frac{1}{\omega + i\eta - \epsilon_p}, \quad (44)$$
where the self-energy correction is ignored [noting $\epsilon_p = p^2/(2m)$]. For the two-body sector, we can rewrite the lowest-order two-body propagator as

$$\Pi_{q,j}^{\text{ret}}(\omega) \equiv \Pi_{q,0}(\omega) + \Xi_{q,j}(\omega),$$

where

$$\Pi_{q,0}(\omega) = \sum_p \frac{1}{\omega + i\eta - \epsilon_{p+q/2} - \epsilon_{-p+q/2}}$$

and

$$\Xi_{q,j}(\omega) = -\sum_p \frac{f_j(\epsilon_{p+q/2}) + f_j(\epsilon_{-p+q/2})}{\omega + i\eta - \epsilon_{p+q/2} - \epsilon_{-p+q/2}}$$

are the in-vacuum two-body Green’s function and the medium correction, respectively (for more details, see e.g., Refs. 33, 35). Taking $\alpha^2 = q^2/4 - m\omega - i\delta$, we can analytically obtain

$$\Pi_{q,0}(\omega) = -\frac{mA}{2\pi^2} + \frac{m\alpha}{2\pi^2} \tan^{-1} \left( \frac{A}{\alpha} \right),$$

where $A$ is an ultraviolet cutoff. Note that $A$ is renormalized via

$$\frac{m}{4\pi a} = \frac{1}{g} + \frac{mA}{2\pi^2},$$

which leads to

$$\frac{1}{\Pi_{q,j}^{\text{ret}}(\omega)} = \frac{m}{4\pi a} - \Pi_{q,j}^{\text{ret}}(\omega) - \frac{mA}{2\pi^2} \sim \frac{m}{4\pi a} - \Xi_{q}(\omega) - \frac{m\alpha}{4\pi}$$

where the ultraviolet divergence is cancelled ($\tan^{-1}(A/\alpha) \approx \pi/2$ is used in the second line).

In the dilute limit, the fermionic medium correction $\Xi_{q,R}(\omega)$ is negligible. In this case, one can approximately obtain

$$G_{q,R}^{\text{ret}}(\omega) \simeq \Pi_{q,0}(\omega) [1 - g\Pi_{q,0}(\omega)]^{-1}.$$  

(51)

where $G_{q,R}^{\text{ret}}(\omega)$ does not involve any poles on the real frequency axis (i.e. bound states) at $a^{-1} < 0$. Note that the two-body continuum exists above $\omega = q^2/(4m)$. In the weak-coupling side ($a < 0$), we obtain

$$B_{q,R}(\omega) = -2 \text{Im} G_{q,R}^{\text{ret}}(\omega) = 0. \quad (\omega < q^2/4m).$$

(52)

Simultaneously, the frequency integration is restricted as $\omega > 0$. This fact indicates that particles in the reservoir L are transferred to the two-body continuum in the reservoir R via the two-body tunneling process in the weak-coupling side ($a < 0$). On the other hand, in the strong-coupling limit ($a \to +\infty$), we obtain 53, 55

$$G_{q,R}^{\text{ret}}(\omega) \simeq \left( \frac{mA}{2\pi^2} \right)^2 \frac{8\pi}{m^2a} \frac{1}{\omega + i\eta - \frac{q^2}{4m} + E_b} \quad (A \to \infty),$$

(53)

which is proportional to the bosonic Green’s function of a bound molecule with the binding energy $E_b = 1/(ma^2)$. Thus, in the strong-coupling regime ($a > 0$), particles in the reservoir L can be transferred to the molecular bound states in the reservoir R via the two-body tunneling process.

**LARGE-BIAS LIMIT**

In the main text, we considered a situation where fermions in the strongly-correlated reservoir L with a finite density $N_L$ go through the tunneling junction to the dilute reservoir R with a vanishing density $N_R \to 0$, i.e., $\mu_R \to -\infty$ (see
Fig. 5 shows the calculated Fano factor $F$ with different tunneling-coupling ratio $\tau_{2,\text{ren.}}/\tau_1$ in the entire BCS-BEC crossover regime at $T_L/T_{F,L} = 0.3$. While in the main text we employed $\tau_{2,\text{ren.}}/\tau_1 = 1$, this ratio depends on the actual detailed setups in each experiment. If the two-body tunneling is relatively strong as $\tau_{2,\text{ren.}}/\tau_1 = 10$, $F$ is close to 2 even in the weak-coupling side [(4) $\zeta = -1$]. However, $F$ decreases at weaker coupling even in this case. On the other hand, in the case with $\tau_{2,\text{ren.}}/\tau_1 = 0.1$, $F$ remains to be close to 1 even around unitarity. Nevertheless, $F$ rapidly increases around $(k_{F,L}a)^{-1} = 0.3$ and consequently reaches $F = 2$ in the strong-coupling limit.

In this way, the detailed structure of the tunneling junction affects how $F$ increases in the BCS-BEC crossover regime. However, our conclusion that $F = 1$ and $F = 2$ are achieved in the BCS and BEC limits, respectively, is unchanged even for different tunneling-coupling ratios. In other words, the pair tunneling process inevitably occurs in the strong-coupling regime even for an infinitesimally small pair-tunneling coupling $\tau_2$. This is a natural consequence in the sense that the system is dominated by bound molecules and hence there are no single-particle states in such a regime.

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