Landauer formulation of photon transport in driven systems

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Understanding the behavior of light in nonequilibrium scenarios underpins much of quantum optics and optical physics. While lasers provide a severe example of a nonequilibrium problem, recent interests in the near-equilibrium physics of so-called photon gases, such as in Bose condensation of light or in attempts to make photonic quantum simulators, suggest one re-examine some near-equilibrium cases. Here we consider how a sinusoidal parametric coupling between two semi-infinite photonic transmission lines leads to the creation and flow of photons between the two lines. Our approach provides a photonic analog to the Landauer transport formula, and using nonequilibrium Green’s functions, we can extend it to the case of an interacting region between two photonic leads where the sinusoid frequency plays the role of a voltage bias. Crucially, we identify both the mathematical framework and the physical regime in which photonic transport is directly analogous to electronic transport and regimes in which other behavior such as two-mode squeezing can emerge.

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I. INTRODUCTION

Quantum systems have dynamics that appear to be beyond the capacity of classical computers to simulate as the size of the system increases. However, a controlled quantum simulator may enable an understanding of such systems that eludes classical description, as the emulation of one system by another can take full advantage of the underlying quantum evolution [1,2]. One promising avenue for quantum simulation uses massless bosons—typical photons—as the constituent particles and examines the phases of matter that can arise with the inclusion of interactions between these particles [3–6]. Perhaps the most dramatic possibilities arise in circuit quantum electrodynamics (QED) [7], where the Josephson effect provides a strong microwave nonlinearity, though similar improvements are now becoming available in semiconductor, molecular, and atomic nonlinearities in small optical domain cavities [8–10]. These photonic systems are particularly interesting given our ability to control the dispersion relation of the particles, including, e.g., the creation of effective mass [11,12] or synthetic gauge fields [13–15] as well as the character of their interaction. As a starting point, Bose-Einstein condensation of photons has been observed in recent experiments using cavity polaritons [11,16–18] or with dye microcavities [12] using these ideas.

Unfortunately, the vacuum is the typical ground state for such systems, and thus efforts for quantum simulation with light have focused on driving systems far from equilibrium to provide sufficient numbers of photons. This makes predicting the dynamics and steady-state behavior an outstanding challenge [19,20]. On the other hand, electronic transport theory, pioneered in the works of Landauer [21–23], Büttiker [24], and Imry [23], has successfully dealt with a different problem: What is the quantum version of Ohm’s law, i.e., the relationship between chemical potential difference (voltage) and particle flux (current), for describing the motion of electrons in mesoscopic systems [21–27]? Of particular use have been mathematical tools such as nonequilibrium Green’s function methods [26–29], which enable predictions for systems even at large voltage bias and with strong interactions.

In this article we consider whether a photonic version of the Landauer-type transport exists and find that for parametrically coupled semi-infinite leads (transmission lines), a natural photonic voltage arises with an associated Ohm’s law-type behavior for the photon flux. Our results rely upon the most recent of several approaches for developing a photonic equivalent to this voltage-bias [30–33], including equilibration of light coupled to electrons flowing in a diode [33–35] and, more recently, parametrically coupled photonic systems [36]. Specifically, we derive the nonequilibrium transport of light under the parametric coupling scheme using nonequilibrium Green’s function (NEGF) formalism. We study the photon flux as the equivalent of a current through a parametrically driven mesoscopic region and show that the photon flux formula can be understood in the Landauer sense, as a transport from a chemical potential imbalance from the parametric coupling, with the addition of an anomalous particle-nonconserving squeezing term. Intuitively, our result connects the photon flow between a low-frequency bath and an optical bath as mediated by a mesoscopic, interacting region. Thus we provide a rigorous framework for studying such near-equilibrium photonic systems without resorting to ad hoc tools for steady-state dynamics. Furthermore, our result predicts a quantitative link between the photon flux and the Green’s function, which provides a possible testing ground for photonic quantum simulations even without particle number conservation.

II. PHOTON TRANSPORT THROUGH A TRIVIAL SCATTERER

We start by developing our photonic analog to voltage bias. Consider a photonic (optical or microwave) system coupled to two baths: one associated with the typical decay

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FIG. 1. (a) Our conceptually simplest system of two semi-infinite leads, with a time-dependent coupling between them and a photodetector connecting to the right lead. (b) A potential physical implementation with a Josephson parametric coupler, driven with a flux bias line, between two transmission lines.

of excitations into other modes via, e.g., imperfect mirrors, while the other is associated with a second bath coupled time dependently with fast sinusoidal variation of the coupling constant at angular frequency $\omega_p$. In particular, in Ref. [36], one of us showed that a time-dependent bath coupling can lead to the equilibrium of a small system best described by a grand canonical ensemble distribution, i.e., a system of photons with a chemical potential. However, in that work crucial questions—such as what happens when coupled to two baths—were largely detailed heuristically. Here we focus on building a formalism, analogous to the finite-bias Green’s function approach for electronic transport. In particular, we describe the two baths as semi-infinite transmission lines for our purposes, with the parametrically coupled bath being the left lead and the natural bath corresponding to photon loss being the right lead, which could correspond to an outgoing optical signal to be measured with a photodetector (Fig. 1). This is now analogous to electronic transport at finite voltage bias, where the voltage is equivalent to the chemical potential $\hbar \omega_p$.

As a toy model, and to help develop the formalism, we start with the simplest setup in which the scatterer is trivial—a flux bias line, between two transmission lines.

We may organize the Hamiltonian in a matrix form

$$H = \frac{1}{2} \tilde{p}^T \tilde{p} + \frac{1}{2} \tilde{u}^T \mathbf{K} \tilde{u}$$

by introducing displacement and momentum vectors

$$\tilde{u} = \begin{pmatrix} \tilde{u}_L \\ \tilde{u}_R \end{pmatrix}, \quad \tilde{p} = \begin{pmatrix} \tilde{p}_L \\ \tilde{p}_R \end{pmatrix},$$

with elements $(u_L(R), a(\beta)) \equiv u_{a(\beta)}$ and $(p_L(R), a(\beta)) \equiv p_{a(\beta)}$. $\tilde{u}$ and $\tilde{p}$ follow the equal time commutation relation

$$[\tilde{u}(t), \tilde{p}(t)] = i \hbar \mathbf{I}.$$ 

Here $\mathbf{I}$ is the identity matrix, and $\mathbf{T}$ denotes the matrix transpose. $\mathbf{K}$ is a symmetric spring’s constant matrix and can be further separated into diagonal and off-diagonal parts

$$\begin{pmatrix} \mathbf{D} & \mathbf{0} \\ \mathbf{0} & \mathbf{D}^R \end{pmatrix}, \quad \mathbf{V}(t) = \begin{pmatrix} 0 & \mathbf{V}^{LR}(t) \\ \mathbf{V}^{RL}(t) & 0 \end{pmatrix}.$$ 

Here $D_{aa'} = \alpha_0^2 \delta_{aa'}, D_{bb'} = \alpha_0^2 \delta_{bb'},$ and $V_{\beta \gamma}^{LR}(t) = V_{\beta \gamma}^{RL}(t) = \cos(\alpha_0 t)\delta_{\beta \gamma}$.

Assume the two subsystems were initially decoupled and in their own thermal equilibrium, and the parametric coupling is adiabatically turned on at $t = -\infty$ and turned off at $t = \infty$. Our goal is to find the photonic current transported between the two ends and express it in a Landauer-like formula in order to predict the current based on an effective chemical potential difference analogous to a voltage bias.

The current on the right at some later time $t$ is defined as the temporal change of the total number of photons in the right transmission line $N_R = \sum_{\beta} b^\dagger_{\beta} b_{\beta}$, which corresponds to an expected photodetector signal. We have

$$J_R(t) \equiv \langle \dot{N}_R(t) \rangle = \left\langle \left. \frac{d}{dt} \sum_{\beta} b^\dagger_{\beta} b_{\beta}(t) \right|_{t'=\omega_p} \right\rangle.$$ 

The angular bracket denotes ensemble average over the initial equilibrium density of states, while the operators are in the Heisenberg picture. According to the Heisenberg equation of motion,

$$\dot{N}_R(t) = -\frac{1}{\hbar} \left\langle \frac{\partial}{\partial t} \left[ \tilde{u}_R(t') \tilde{V}_{RL}(t') \tilde{u}_L(t) \right] \right|_{t'=\omega_p},$$

where $\tilde{V}_{\beta \gamma}^{RL}(t) \equiv V_{\beta \gamma}^{RL}(t)/\omega_{\beta}$.

One can connect the current expression with Keldysh Green’s functions [28] by introducing the nonequilibrium lesser Green’s function defined as

$$G^{<}(t, t') \equiv -\frac{i}{\hbar} \langle \tilde{u}(t')\tilde{u}^T(t) \rangle,$$

which can also be split into four blocks associated with left and right transmission lines. We can write the current using the lesser Green’s function as

$$J_R(t) = -i \left\langle \frac{\partial}{\partial t} \text{Tr} [G^{<}_{LR}(t, t') \tilde{V}_{RL}(t')] \right|_{t'=\omega_p}.$$ 

The trace here means tracing over photon states $\alpha$. We now follow the standard Keldysh formalism (NEGF formalism) [28] to study the transport formula [25–27,37]. Since we define our Green’s functions on displacement operators $u$ instead of photon creation operators $a^\dagger$, our
problem structurally resembles more the thermal transport cases [37] than electronic ones. We remind the reader here that since the parametric coupling varies with time and allows pair production and annihilation mechanisms, many identities and tricks in previous works involving steady-state or particle-conserving assumptions cannot be applied here.

The equation of motion of the contour ordered Green’s function defined on the Keldysh contour C follows

$$\frac{\partial^2}{\partial t^2} G^r(r, r') + K G^r(r, r') = -\delta(r, r') I,$$

while the noninteracting equilibrium Green’s function $g^r(r, r')$ follows the equation of motion

$$\frac{\partial^2}{\partial t^2} g^r(r, r') + D g^r(r, r') = -\delta(r, r') I.$$  

One can easily verify that $G^r(r, r')$ follows the Dyson equation

$$G^r(r, r') = g^r(r, r') + \int_C d\tau'' g^r(r, \tau'') V(\tau'') G^r(\tau'', r').$$

Using the Langreth theorem of analytic continuation [29], the lesser Green’s function can be expressed as an integral form on the real axis:

$$G_{L,R}^r(t, t') \approx \int_{-\infty}^{\infty} dt_1 \left[ g^r_{L}(t_1, t) V^{L}(t_1) g^r_{R}(t_1, t') + g^r_{L}(t_1, t) V^{R}(t_1) g^r_{R}(t_1, t') \right] + \mathcal{O}(\lambda^2).$$

Here the $r$ and $a$ superscripts stand for retarded and advanced Green’s functions, and we treat $\lambda$ as a perturbation. The equilibrium Green’s functions used in the $G_{L,R}^r(t, t')$ expression are given by

$$(g_{L}^r)_{\alpha}(t, t_1) = -\frac{i}{2\omega_{\beta}} \theta(t - t_1) e^{-i\omega_{\beta}(t - t_1)} - e^{i\omega_{\beta}(t - t_1)},$$

$$(g_{R}^r)_{\beta}(t_1, t') = -\frac{i}{2\omega_{\alpha}} \{ n_{R}(\epsilon_{\beta}) e^{-i\omega_{\alpha}(t_1 - t')} + [1 + n_{R}(\epsilon_{\beta})] e^{i\omega_{\alpha}(t_1 - t')},$$

$$(g_{L}^a)(t_1, t) = -\frac{i}{2\omega_{\alpha}} [ n_{L}(\epsilon_{\alpha}) e^{-i\omega_{\alpha}(t_1 - t)} + [1 + n_{L}(\epsilon_{\alpha})] e^{i\omega_{\alpha}(t_1 - t)}],$$

$$(g_{R}^a)(t_1, t') = -\frac{i}{2\omega_{\beta}} \theta(t_1 - t') [ e^{i\omega_{\beta}(t_1 - t')} - e^{-i\omega_{\beta}(t_1 - t')}].$$

Here $n_{L,R}(\epsilon_{\alpha,\beta}) = (e^{\epsilon_{\alpha,\beta} - \mu_{L,R}}/\kappa_{B}T - 1)^{-1}$ are the bosonic occupation number in left and right transmission lines. The chemical potentials $\mu_L = \mu_R = 0$ for photons, $\kappa_B$ is the Boltzmann constant, and $T$ is the initial temperature of the system.

By inserting the expression for $G_{L,R}^r(t, t')$, the current is now

$$J_R(t) = -i \left( \frac{\partial}{\partial t} \right) \left[ \int_{-\infty}^{\infty} dt_1 \left[ g^r_{L}(t_1, t) V^{L}(t_1) g^r_{R}(t_1, t') + g^r_{L}(t_1, t) V^{R}(t_1) g^r_{R}(t_1, t') \right] V^{R}(t_1) \right] \bigg|_{t' = t}.$$

$$= -i \left( \frac{\partial}{\partial t} \right) \sum_{\alpha, \beta} \int_{-\infty}^{\infty} dt_1 \lambda_{\alpha, \beta}^{2} \cos(\omega_{\alpha}t_1) \cos(\omega_{\alpha}t') \left[ (g_{L}^r)^{\alpha}(t_1, t) (g_{R}^r)^{\beta}(t_1, t') + (g_{L}^r)^{\alpha}(t_1, t) (g_{R}^a)^{\beta}(t_1, t') \right] \right] \bigg|_{t' = t},$$

$$= \sum_{\alpha, \beta} \lambda_{\alpha, \beta}^{2} \int_{0}^{\infty} d\tau \left[ \cos(\omega_{\alpha}t) + \cos(\omega_{\alpha}(2t - \tau)) \right] \left[ \cos(\omega_{\beta} - \omega_{\alpha}) \tau \right] [ n_{L}(\epsilon_{\alpha}) + n_{R}(\epsilon_{\beta}) ]$$

$$+ \cos(\omega_{\beta}t) + \cos(\omega_{\beta}(2t - \tau)) \right] \cos(\omega_{\beta} + \omega_{\alpha}) \tau \left[ n_{L}(\epsilon_{\alpha}) + n_{R}(\epsilon_{\beta}) + 1 \right].$$

In the last equality we have used the identity $\cos(\omega_{\alpha}t_1) \cos(\omega_{\alpha}t') = \{ \cos(\omega_{\alpha}t) + \cos(\omega_{\alpha}(2t - \tau)) \}/2$ and changed the integral variable to $\tau = t - t_1$. The only explicit $\tau$ dependence arises in the $\cos(\omega_{\alpha}(2t - \tau))$ factor. Averaging over one pump cycle $2\pi / \omega_{\alpha}$ takes this factor to zero. We thus neglect those terms with $\cos(\omega_{\alpha}(2t - \tau))$ in the spirit of the rotating wave approximation.

We assume the coupling constant only depends on the mode energy, $\lambda_{\alpha, \beta} = \lambda(\epsilon_{\alpha}, \epsilon_{\beta})$, and take the continuum limit of energy so that

$$\sum_{\alpha, \beta} \int_{0}^{\infty} d\epsilon_{\alpha} d\epsilon_{\beta} \rho_{L}(\epsilon_{\alpha}) \rho_{R}(\epsilon_{\beta}).$$

Here $\rho_{L}$ and $\rho_{R}$ are the energy density of states in the left and right transmission lines. Note that $\int_{0}^{\infty} \cos(\omega_{\alpha}t_1) dt_1 = \pi \delta(\omega_{\alpha} - \omega_{\alpha})$. We now arrive at a current formula with three terms:

$$J_R = \int_{-\hbar\omega_{p}}^{\infty} d\epsilon T(\epsilon, -\hbar\omega_{p}) [ n_{L}(\epsilon) - n_{R}(\epsilon - \hbar\omega_{p}) ] + \int_{-\hbar\omega_{p}}^{\infty} d\epsilon T(\epsilon - \hbar\omega_{p}, \epsilon) [ n_{L}(\epsilon - \hbar\omega_{p}) - n_{R}(\epsilon) ]$$

$$+ \int_{0}^{\hbar\omega_{p}} d\epsilon T(\epsilon, \hbar\omega_{p} - \epsilon) [ n_{L}(\epsilon + \hbar\omega_{p} - \epsilon) + n_{R}(\epsilon) ] + 1].$$
and as the IR divergence is approached for lower dimensional systems, an appropriate inclusion of pump depletion will be necessary to develop a complete understanding of the problem.

The first line of Eq. (16) can be interpreted as a Landauer-like transport with an effective chemical potential $\hbar \omega_p$ on the right transmission line, and the second line represents a Landauer-like transport with an effective chemical potential $\hbar \omega_p$ on the left. The third line is a particle-nonconserving term due to pair creation and annihilation mechanisms allowed by the oscillating $u$-$u$ type coupling. Two-mode squeezed states of light [38] are generated through this mechanism with the photon pairs entangled. One will expect a thermal state when tracing over the output modes on one side of such photon pairs.

Note that the current formula is consistent with Fermi’s golden rule: The parametric coupling $\cos(\omega_p t)$ only allows transition with $E_f - E_i = \Delta E = \pm \hbar \omega_p$, where $E_f$ and $E_i$ are the energies of the final and initial states. For $\omega_p = 0$, the current equation reduces to the usual Landauer form proportional to $n_L(\epsilon) - n_R(\epsilon)$, which is essentially zero when the two transmission lines are at the same temperature.

The particle-nonconserving nature of the problem is manifested by identifying the anomalous current $\bar{J}_A = (\bar{J}_R + \bar{J}_L)/2 = \int_0^{\hbar \omega_p} d\epsilon T(\epsilon, \hbar \omega_p - \epsilon)[n_L(\epsilon) + n_R(\hbar \omega_p - \epsilon) + 1]$, which is only zero when $T(\epsilon, \hbar \omega_p - \epsilon) = 0$ throughout the range, as is the case in Fig. 2(b). This term can also be understood in Fermi’s golden rule point of view, considering the harmonic perturbation $H_F(t)$. According to Fermi’s golden rule, the pair creation (annihilation) rates $R_c$ ($R_a$) are

$$R_c = \int_0^{\hbar \omega_p} d\epsilon T(\epsilon, \hbar \omega_p - \epsilon)[n_L(\epsilon) + 1][n_R(\hbar \omega_p - \epsilon) + 1],$$

$$R_a = \int_0^{\hbar \omega_p} d\epsilon T(\epsilon, \hbar \omega_p - \epsilon)n_L(\epsilon)n_R(\hbar \omega_p - \epsilon).$$

The net creation rate is thus $R_c - R_a = \bar{J}_A$.

One can find the nonequilibrium transport part of the current by subtracting the anomalous squeezing (particle-nonconversing) term $\bar{J}_A$, and we are left with the normal current $\bar{J}_N = (\bar{J}_R - \bar{J}_L)/2$, the first two lines of Eq. (16). We note here that the asymmetry between right and left is necessary for the transport to occur; the first two lines of Eq. (16) will cancel each other otherwise.

To focus on the transport mechanism only, we consider an energy gap on the right transmission line [see Fig. 2(b)] such that $\forall \alpha, \beta, \epsilon_\beta > \hbar \omega_p, \epsilon_\alpha > \epsilon_\alpha$. This gap setup prevents the pair creation and annihilation mechanisms, leaves us with a conserved current, and permits a direct photonic analog to electronic transport. The system follows the transport formula $\bar{J}_N = \int_{\epsilon_{\alpha,\min}}^{\epsilon_{\alpha,\max} + \hbar \omega_p} d\epsilon T(\epsilon - \hbar \omega_p, \epsilon)[n_L(\epsilon - \hbar \omega_p) - n_R(\epsilon)]$, which is equivalent to a nonequilibrium transport current under a chemical potential imbalance $\mu_L = \hbar \omega_p, \mu_R = 0$.

One can see the resemblance between our gapped transport equation and the I-V (current-voltage) characteristic of an ideal light-emitting diode (LED) [33] by relating the chemical potential $\hbar \omega_p$ to $qV$ and the gap energy $\epsilon_\beta_{\min}$ to the photon energy threshold $\epsilon_\beta$ and working under the region $\epsilon_{\alpha,\max} + \hbar \omega_p \gg k_B T$. However, we cannot yet make a direct connection mathematically with the somewhat different problem of electron transport through a diode combined with emission of photon into an interacting region.

III. PHOTON TRANSPORT THROUGH A MESOSCOPIC CENTRAL REGION

Now we consider a more generic case with a center mesoscopic region placed between the transmission lines, with parametric coupling between the center region and the left transmission line (see Fig. 2). We replace the time-dependent barrier $H_F(t)$ with $H_C + H_{CL}(t) + H_{CR}$, and the Hamiltonian becomes

$$H = H_L + H_C + H_R + H_{CL}(t) + H_{CR},$$

$$H_L = \sum_\alpha \epsilon_\alpha a_\alpha^\dagger a_\alpha, \quad H_R = \sum_\beta \epsilon_\beta b_\beta^\dagger b_\beta,$$

$$H_C = \sum_{\gamma,\gamma'} t_{\gamma\gamma'} c_\gamma^\dagger c_{\gamma'} + H_{\text{int}},$$
\[ H_{\text{CL}}(t) = \cos(\alpha t) \sum_{a,\gamma} \lambda_{a\gamma} u_a \bar{u}_\gamma, \]
\[ H_{\text{CR}} = \sum_{\beta,\gamma} \bar{\lambda}_{\beta\gamma} u_\beta \bar{u}_\gamma. \]

The summation index \( \gamma \) labels the states in the center with energies \( \hbar\omega_{\gamma} = \epsilon_{\gamma} \) and photon annihilation operators \( c_{\gamma} \). Note that we place all time dependence in the center to left coupling \( H_{\text{CL}}(t) \). The central region can contain some nonlinear interacting term \( H_{\text{int}} \) as well as nontrivial single-particle potential effects from \( t_{\gamma\gamma} \). We again assume the subsystems were initially in their own thermal equilibrium before the parametric coupling adiabatically turned on at \( t = -\infty \).

The current on the right can be expressed as
\[ J_R(t) = -i \int_{t_0}^{t} dt' \langle \text{Tr}[G_{\text{CR}}^{-1}(t,t')V^R] \rangle_{t' = t}. \]

Note that for the case of interacting center, there will be additional contribution to the self-energy depending on the details of \( H_{\text{int}} \).

Specifically, the center greater and lesser Green’s function follows
\[ G_{\text{CC}}^\pm(t,t') = \int dt_1 dt_2 G_{\text{CC},\pm}(t,t_1) \Sigma_{\text{tot}}^\pm(t_1,t_2) G_{\text{CC}}^\pm(t_2,t'), \]
which allows us to further simplify the current expression.

The center Green’s function under parametric coupling can be expanded with the harmonics of the coupling frequency \( \omega_p \). Specifically, \( G_{\text{CC},\pm}(t,\omega) = \sum_n G_{\text{CC},\pm}(\omega) e^{i n \omega_p t}, n \in Z \). Under the rotating wave approximation, we neglect fast oscillating terms with \( n \neq 0 \) and keep only the \( n = 0 \) steady part of the current. The time averaged current is now
\[ \overline{J}_R = \int_{-\infty}^{\infty} d\omega \frac{d\omega}{2\pi} \text{Tr}[G_{\text{CC},+}(\omega) \Sigma_{\text{R}}^+(\omega) + G_{\text{CC},-}(\omega) \Sigma_{\text{R}}^-(\omega)]. \]

Since the current is real, \( \overline{J}_R = (\overline{J}_R + \overline{J}_R^*)/2 \). Using the general identity \( G^{-} - G^{+} = G^{R} - G^{\Lambda} \) and identities for steady-state Green’s functions in the frequency domain \( [G^{R}(\omega)]^1 = G^{R}(\omega), [G^{\Lambda}(\omega)]^1 = -G^{R}(\omega), \)
\[ \overline{J}_R = \int_{-\infty}^{\infty} d\omega \frac{d\omega}{2\pi} \text{Tr}[\{G_{\text{CC},+}(\omega) - G_{\text{CC},-}(\omega)\} \Sigma_{\text{R}}^+(\omega) \]
\[ - G_{\text{CC},-}(\omega) \{ \Sigma_{\text{R}}^+(\omega) - \Sigma_{\text{R}}^-(\omega) \}]. \]

Expanding the noninteracting center greater and lesser Green’s functions to the leading order term yields
\[ G_{\text{CC},+}(t,t') \approx \int dt_1 dt_2 G_{\text{CC},+}(t,t_1) \Sigma_{\text{tot}}(t_1,t_2) g_{\text{CC}}(t_2,t') + O(\lambda^2), \]
\[ G_{\text{CC},-}(\omega) \approx g_{\text{CC},-}(\omega) \Sigma_{\text{tot},-}(\omega) g_{\text{CC}}(\omega) \]
\[ = g_{\text{CC}}(\omega) \left[ V^R g_{\text{CC}}(\omega) V^C \right. \]
\[ + \frac{1}{4} \frac{V^L g_{\text{CC}}(\omega) + \omega_p}{V^C} \]
\[ + \frac{1}{4} \frac{V^L g_{\text{CC}}(\omega) - \omega_p}{V^C} \]
\[ g_{\text{CC}}(\omega). \]

By inserting the equilibrium Green’s function for left and right transmission lines \( g_L \) and \( g_R \), we arrive at a formula similar to the trivial scatterer problem:
\[ \overline{J}_R = \int_{h\omega_p}^{\infty} \frac{d\epsilon}{\hbar^3} T_C(\epsilon, \epsilon - h\omega_p) \left[ n_L(\epsilon) - n_R(\epsilon - h\omega_p) \right] \]
\[ + \int_{h\omega_p}^{\infty} \frac{d\epsilon}{\hbar^3} T_C(\epsilon - h\omega_p, \epsilon) \left[ n_L(\epsilon - h\omega_p) - n_R(\epsilon) \right] \]
\[ + \int_{0}^{\hbar\omega_p} d\epsilon T_C(\epsilon, h\omega_p - \epsilon) \left[ n_L(\epsilon) + n_R(\hbar\omega_p - \epsilon) + 1 \right], \]
where the center transmission function is
\[ T_C(\epsilon_a, \epsilon_b) = \frac{\pi h^3}{8} \frac{1}{\text{Tr}[g(\epsilon_b) \Lambda(\epsilon_b) g(\epsilon_b) \Lambda(\epsilon_a) \Lambda(\epsilon_a)]}, \]
\[ \Lambda(\epsilon_a) \gamma_{1,2} = \rho_L(\epsilon_a) \lambda_{\gamma_{1,2}}(\epsilon_a) / \epsilon_a, \]
\[ \Lambda(\epsilon_b) \gamma_{1,2} = \rho_R(\epsilon_b) \lambda_{\gamma_{1,2}}(\epsilon_b) / \epsilon_b. \]
We again have the first two lines of Eq. (29) as the Landauer-like transport terms, and the last line is the particle-nonconserving part due to the oscillating $u-u$ type coupling. The system will undergo nonequilibrium transport with an effective chemical potential imbalance $\hbar \omega_p$ under specific gap setups, and the current expression resembles the I-V characteristic of an ideal light-emitting diode.

IV. CONCLUSIONS

We have derived the photonic flux between different baths parametrically coupled to an intermediate system and found a Landauer-like transport formula for noninteracting centers. However, we also have another regime, with a particle-nonconserving term, which we can interpret as a two-mode squeezing output. The consequences of this latter regime for observation and even application remain to be explored and are beyond the scope of the present work. We have also shown a potential extension of these techniques at the formal level to the interacting case, but suggest that applying these results, e.g., to photon-blockaded systems to see the nonclassical light output would be an intriguing next step.

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