Cavity-Enhanced Transport of Excitons

Johannes Schachenmayer,[1] Claudiu Genes,[2] Edoardo Tignone,[3] and Guido Pupillo[3]
[1]JILA, NIST, Department of Physics, University of Colorado, 440 UCB, Boulder, Colorado 80309, USA
[2]Institut für Theoretische Physik, Universität Innsbruck, Technikerstrasse 25, A-6020 Innsbruck, Austria
[3]IPCMS (UMR 7504) and ISIS (UMR 7006), Université de Strasbourg and CNRS, Strasbourg 67000, France
(Received 8 September 2014; published 12 May 2015)

We show that exciton-type transport in certain materials can be dramatically modified by their inclusion in an optical cavity: the modification of the electromagnetic vacuum mode structure introduced by the cavity leads to transport via delocalized polariton modes rather than through tunneling processes in the material itself. This can help overcome exponential suppression of transmission properties as a function of the system size in the case of disorder and other imperfections. We exemplify massive improvement of transmission for excitonic wave packets through a cavity, as well as enhancement of steady-state exciton currents under incoherent pumping. These results may have implications for experiments of exciton transport in disordered organic materials. We propose that the basic phenomena can be observed in quantum simulators made of Rydberg atoms, cold molecules in optical lattices, as well as in experiments with trapped ions.

DOI: 10.1103/PhysRevLett.114.196403 PACS numbers: 71.35.-y, 05.60.Gg, 37.30.+i, 81.05.Fb

Understanding the transport properties of quanta and correlations and how to make this transport efficient over large distances are questions of fundamental importance in a variety of fields, ranging from experiments with cold atoms and ions [1–4], to quantum information theory [5–7], to (organic) semiconductor and solar cell physics [8–10]. In most realistic situations, transport efficiency is known to be strongly inhibited by disorder. For example, Anderson-type localization of single-particle eigenstates [11] in disordered media implies an *exponential* suppression of transmission; i.e., over a distance of N sites it decays as $T \propto \exp(-N)$. In this work we show how in general exponential suppression of energy transport via atomic and molecular excitons can be overcome by coupling the excitons to the structured vacuum field of a Fabry-Perot cavity placed transverse to the propagation direction. In one dimension (1D), this trades the exponential suppression for a decay which is at most *algebraic*, $T \propto N^{-2}$, a massive enhancement that should be observable for realistic exciton-cavity couplings, system sizes, disorder strengths, and even at room temperature [12–22]. While here we focus on *exciton* transport, our work was originally inspired by the first breakthrough experiments on *charge* transport in molecular semiconductors in the strong-coupling regime [12]. In principle, the observed effect may open the way towards utilizing molecular materials as inexpensive and flexible alternatives to traditional silicon-based semiconductors [8,23–29].

Here, we provide a theoretical understanding of enhanced exciton transport for a model of two-level systems embedded in a cavity in the limit of strong collective light-exciton coupling. We note that in these systems, strong collective coupling has been already demonstrated, and even used, e.g., to modify intrinsic material properties such as the work function [30]. On the other hand, our model also applies to artificial media such as cavity-embedded Rydberg lattice gases [31,32], polar molecules in optical lattices [33,34], or ions in linear Paul traps [35,36]. In these systems, large couplings [37] and reduced decoherence from spontaneous emission may allow for demonstrating essentially instantaneous coherent transport of excitonic wave packets over large distances with close-to-unit efficiency, $T \propto 1$. These experiments can analyze transport in systems with many excitations and in high dimensions, where modern numerical methods become inefficient [38–40]. This may contribute to improving our understanding both of transport in real materials and of fundamental properties of information transport in strongly correlated light-coupled systems [1,2,41,42].

The *model* we consider consists of a chain of N two-level systems or "spins" with local states $|\uparrow\rangle_i$ and $|\downarrow\rangle_i$ that are embedded in a cavity. The coupling to the cavity is governed by the Tavis-Cummings Hamiltonian $H_{\mathrm{cav}} = g\sum_i(\sigma_i^+ a + \sigma_i^- a^\dagger)$, with g the coupling strength, a ($a^\dagger$) the destruction (creation) operator for the cavity photon, and $\sigma_i^\pm$ the Pauli spin raising or lowering operators for the spin at site i. We restrict our discussion to single excitations in the system. Such a localized excitation (i.e., a state $|\uparrow\rangle_i$) has an energy ω_i ($\hbar \equiv 1$) and can tunnel between neighboring sites, as described by the Hamiltonian $H_0 = \sum_i[\omega_i \sigma_i^+\sigma_i^- - J_i(\sigma_i^+\sigma_{i+1}^- + \sigma_i^-\sigma_{i+1}^+)]$. The tunneling rates J_i can be site dependent, and we define $J_i = J + \delta J_i$, where δJ_i denotes random disorder drawn from a normal distribution with standard deviation δJ. We note that in realistic physical setups the tunneling is typically induced by dipolar long-range forces that give rise to additional long-range hopping terms. These terms are not capable of lifting the exponential suppression of transmission [43,44]. Thus, for simplicity we consider the nearest neighbor tunneling model here. In addition, a coupling between excitons and phonons can give rise to nonlinear terms causing self-trapping effects [48,49]. It can be shown that these terms are very small for realistic parameters and we will neglect them here.

The general dynamics of our system is governed by the master equation $\dot{\rho} = -i[H,\rho] + \sum_\alpha \mathcal{L}_\alpha(\rho)$, with ρ the density matrix and $H = H_0 + H_{\text{cav}}$. The terms $\mathcal{L}_\alpha(\rho) \equiv -\{L_\alpha^\dagger L_\alpha, \rho\} + 2L_\alpha \rho L_\alpha^\dagger$ incorporate all dissipative processes via ordinary Lindblad operators L_α. We consider cavity decay ($L_\kappa \equiv \sqrt{\kappa/2}a$) as well as spontaneous emissions of each spin ($L_{\text{sp em},i} \equiv \sqrt{\gamma_{\text{sp em}}/2}\sigma_i^-$) or dephasing ($L_{\text{deph},i} \equiv \sqrt{\gamma_{\text{deph}}/2}\sigma_i^+\sigma_i^-$), deriving, e.g., from radiative decay and fluctuations in level spacing (vibrations) due to the system being at finite temperature. In the homogeneous situation with $\omega_i = \omega_0$ and $J_i = 0$, H_{cav} is responsible for the formation of dressed modes of the cavity photons and of the collective Dicke states $\sigma_0^\pm \equiv \sum_j \sigma_j^\pm/\sqrt{N}$, named as upper and lower polaritons for $u^\dagger \equiv (a^\dagger + \sigma_0^+)/\sqrt{2}$ and $d^\dagger \equiv (a^\dagger - \sigma_0^+)/\sqrt{2}$, respectively, with energy $\Omega_{u,d} = \omega_0 \pm g\sqrt{N}$. In this work we study two possible scenarios to observe enhancement of exciton transport by exploiting these states: (i) a wave-packet scattering experiment, and (ii) steady state exciton currents under incoherent pumping.

Case (i) is sketched in Fig. 1(a): In addition to the N spins in the cavity, M spins are added to the left and right of the cavity ($\mathcal{N} = 2M + N$), coupled via H_0. We consider a homogeneous level spacing inside and outside of the cavity, with $\omega_i = \omega_0$ for $i = M+1, \ldots, M+N$ and $\omega_i = \omega$

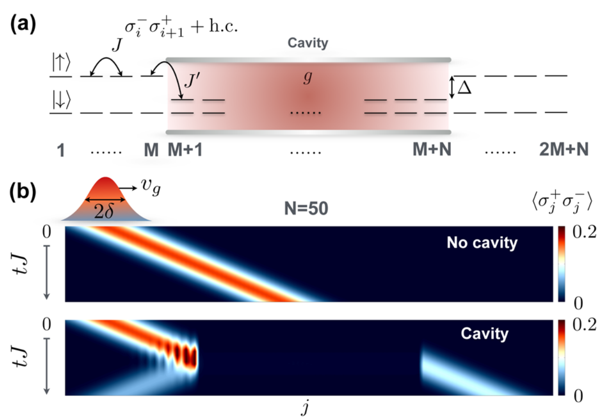

FIG. 1 (color online). Exciton transmission model. Scheme of a chain of coupled two-level systems (tunneling rate J) in which an exciton wave packet propagates from the left into a cavity (group velocity v_g) that is coupled to N spins (cavity-spin coupling g). Under the right conditions, a large portion of the wave packet can be almost instantaneously transmitted to the right side on a time scale $t \ll N/v_g$ [example in panel (b) with $N = 50$, $v_g = 2J$, $\delta = 5$, $\Delta = 69J$, $J' = 10J$, $g = 10J$].

otherwise, and define $\Delta = \omega - \omega_0$. We further denote $J_i = J'$ for $i = M+1$ and $i = M+N$, i.e., at the entrance and exit of the cavity, to allow for impedance effects. At time $t = 0$, a wave packet of excitons, $|\psi(t=0)\rangle \propto \sum_{j=1}^{\mathcal{N}} e^{-iq_0 j} e^{-(j-j_0)^2/(4\delta^2)}|j\rangle$, with width δ (standard deviation) and initial quasimomentum q_0 is injected on the left. Here, $|j\rangle \equiv |\uparrow\rangle_j \bigotimes_{i\neq j} |\downarrow\rangle_i$ denotes the state of a single excitation at site j. The initial displacement from the cavity is $\delta_x = M - j_0$. As an example, we choose $\delta_x = 20$, $\delta = 5$ and $q_0 = \pi/2$ [corresponding group velocity $v_g = 2J\sin(q_0) = 2J$]. We are interested in the wave-packet fraction that for properly tuned parameters can be transferred nearly instantaneously to the right side of the cavity [cf. Fig. 1(b)].

Case (ii), in contrast, concerns a system with sites $i = 1, \ldots, N$ embedded in the cavity. Excitations are incoherently pumped to site $i = 1$ from the left and removed from site $i = N$. This can be achieved via dissipative terms with $L_P \equiv \sqrt{\gamma_P/2}\sigma_1^+$ and $L_{\text{out}} \equiv \sqrt{\gamma_{\text{out}}/2}\sigma_N^-$, respectively. Under these conditions, we calculate the output exciton current, $I_{\text{out}} = \text{tr}[n_N \mathcal{L}_{\text{out}}(\rho)]$ (with $n_i = \sigma_i^+\sigma_i^-$) in the steady state. Similar to the case of Ref. [50], this current arises naturally from the continuity equation $d\langle N_e\rangle/dt = 0 = \text{tr}[N_e d\rho/dt]$, where $N_e = \sum_i n_i$. In the second part of this Letter we show how I_{out} can be dramatically enhanced in the presence of the cavity.

Wave-packet scattering.—In case (i), we first simplify the dynamics by neglecting dissipative terms and disorder (a valid approximation for, e.g., a Rydberg lattice gas [44]). Under these conditions, for $g = 0$ the wave packet ($v_g = 2J$) reaches the right side of the cavity on a *long* time scale $t_l J = \delta_x + 2\delta + N/2$. This corresponds to the time required to hop over N sites plus the time needed to enter and exit the cavity within the light cone. Here we propose using the polariton mode to tunnel N sites almost instantaneously.

The time scale for a single excitation to couple in and out of such a mode is proportional to $\sqrt{N}/g$, and can be exceedingly small for large g. Then, transmission to the right side beyond the free-evolution light cone is possible on an *ultrashort* scale $t_s J = \delta_x + 2\delta \ll t_l J$, limited only by the entrance time in the cavity. The dynamics can then be described via elastic scattering through the cavity, with a quasimomentum dependent transmission function $T_q = |t_q|^2$, and t_q the coefficient appearing in the associated Lippmann-Schwinger equation [51].

The time-independent function T_q determines the transmission properties of the material [52–55], and can be computed exactly for our model. The coefficient has the general form $t_q = -2i\beta/[1 + 2i\beta]$, with $\beta = [2NJ\sin(q)]^{-1}\sum_n |J'|^2/[\omega - 2J\cos(q) - \tilde{\Omega}_n]$. Here, $\tilde{\Omega}_n$ is the nth eigenvalue of the reduced Hamiltonian for the cavity-coupled central N sites of the chain [44]. The resulting T_q in general presents three regions of ballistic transmission (i.e., $T_q = 1$). These correspond to (a) ordinary exciton hopping for $\Delta \sim 0$, with an approximate width $4J$, as well as (b) two peaks for $\Delta \sim \pm g\sqrt{N} - J$. The latter correspond to polariton-mediated transmission, and have an approximate Lorentzian shape with an N-dependent full width at half maximum (FWHM) $w = J'^2/(N|v_g|)$. For large enough strength of the collective exciton-cavity coupling $g\sqrt{N} > \max[w, 4J, \kappa]$ all peaks are well separated, which defines the collective strong coupling regime. In the following we focus on this regime, where in the vicinity of the polariton peaks T_q is found to simplify to

$$T_q = \{1 + N^2 J^2 \sin^2(q)(\omega + J[1 - 2\cos(q)] - \Omega_{u,d})^2/J'^4\}^{-1}. \quad (1)$$

Time-dependent wave-packet scattering can be investigated via numerical exact diagonalization. We define a time-dependent transmission as $T_{t'} = \sum_{j>M+N} \langle \sigma_j^+ \sigma_j^- \rangle_{t'}$, which measures the total number of excitations that reach the right side of the system at a given time t'.

Our goal is to realize large ultrafast transmission via the polariton peaks, i.e., $T_{t'} \sim 1$ at $t' = t_s$. Two conditions have to be met: (i) The detuning Δ has to match the energy of one of the polariton peaks; and (ii) the wave packet has to be sufficiently sharp in quasimomentum space to fit into the energy window w, implying a real-space width on the order of the cavity length. While this can be generally difficult to realize, we find that condition (ii) can be satisfied by a choice $J' \propto \tilde{J}_N \equiv (2\ln 2)^{1/4}\sqrt{N/2\delta}J$ (ensuring an N-independent width), similar to an impedance effect.

In Fig. 2(a) we compare $T_{t'}$ for different Δ, for $t' = t_s$ (red continuous line) and $t' = t_l$ (black dashed line). We choose $N = 100$, large $g = 50J$, and set $J' = 4\tilde{J}_N$. As expected, we find the existence of two distinct polariton peaks, suitable for ballistic transmission on the ultrafast scale t_s. The position and width of the peaks are in agreement with the analytical time-independent predictions of Eq. (1). The peak at $\Delta \sim 0$ instead reflects regular exciton hopping on a time scale $t_l \gg t_s$. Note that here $T_{t_l} < 1$ due to backscattering at the cavity entrance where $J' > J$.

When decreasing the coupling strength g, the exciton dynamics through the cavity slows down considerably [44]. The scattering becomes generally inelastic within t_s: part of the wave-packet energy remains in the cavity and $T_{t_s} < 1$. However, we find that even for moderate couplings, a large fraction of the exciton wave packet is transmitted within t_s. This is shown in Fig. 2(b). There, $\max_\Delta(T_{t_s})$ (i.e., the best achievable T_{t_s} for Δ chosen close to the upper polariton energy) is plotted as a function of g and N: For increasing N, T_{t_s} remains large and constant for a choice $g \sim \sqrt{N}J$. In addition, Fig. 2(c) shows that T_{t_s} vs $g/\sqrt{N}J$ displays a

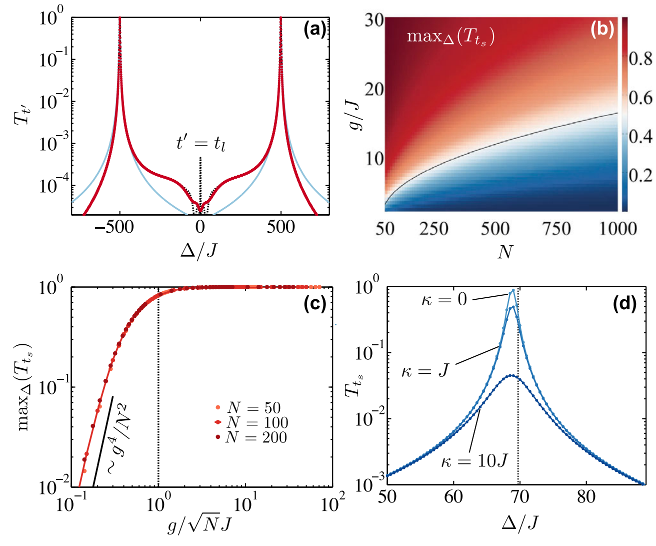

FIG. 2 (color online). Ultrafast transmission of a wave packet with $\delta = 5$, $\delta_x = 20$, and $v_g = 2J$. We choose $J' = 4\tilde{J}_N$ (see text). (a) Long-time and ultrashort transmission (T_{t_l} and T_{t_s}) as function of Δ. The cavity embeds $N = 100$ sites and $g = 50J$ (strong collective coupling regime). Clearly, two peaks of T_{t_s} and T_{t_l} appear at the polariton energies. The numerical calculation (red line) agrees with the analytical result (blue line). T_{t_l} (black dotted line) contains a small $\Delta \sim 0$ peak, corresponding to free evolution. (b) $\max_\Delta(T_{t_s})$ as function of g and N. To keep the ultrafast transmission fixed, $g \propto \sqrt{N}$ (solid line) is required. (c) Crossover into the regime of large ultrafast transmission around $g \sim \sqrt{N}J$ (dotted line). $\max_\Delta(T_{t_s})$ for $N = 50, 100, 200$ is shown as function of $g/\sqrt{N}$ (on top of each other). For small g, $T_{t_s} \sim g^4/N^2$ (solid line). (d) Shrinkage and broadening of transmission peaks for finite cavity decay κ ($N = M = 50$, $g = 10J$).

universal behavior for different N. Here, T_{t_s} reaches large values $\sim 80\%$ for $g = \sqrt{N}J$ improving to 100% when increasing $g/\sqrt{N}J$. This is expected, since for $g > \sqrt{N}J$, we enter the elastic scattering regime, in which the time scale for coupling in and out of the polariton mode becomes negligible. Then, $T_{t_s} \sim 1$ is possible over arbitrarily large distances, if a coupling strength $g \gtrsim \sqrt{N}J$ can be engineered. Interestingly, even for $g \ll \sqrt{N}J$ a significant part of the wave packet is transmitted within t_s. In this regime (inelastic scattering and collective strong coupling), we find a general scaling of $T_{t_s} \sim g^4$ and $T_{t_s} \sim 1/N^2$. Thus, cavity-mediated transmission decreases only *algebraically* with N, which can be important, e.g., when competing against exponential suppression due to disorder.

A lossy cavity ($\kappa \neq 0$) generally decreases T_{t_s} because of loss of exciton population, while the FWHM increases accordingly. Figure 2(d) demonstrates that ultrafast transmission of a large wave-packet fraction is still possible for $\kappa \sim J$ (as, e.g., in a polar molecule setup). We also note that for $\kappa \gg 1$ (after an adiabatic elimination of the cavity mode) dynamics can be described by an all-to-all

Hamiltonian $H_{\rm eff} \approx (2g^2/\kappa)\sum_{i,j}\sigma_i^-\sigma_j^+$. Similar as with $H_{\rm cav}$ above, we find that $H_{\rm eff}$ can give rise to ultrafast transmission. We propose that such a situation could, for example, be observed in experiments with trapped ions, where these type of very long-ranged interactions arise naturally even in the absence of a cavity [44].

In realistic organic semiconductors, disorder is key both in the spatial distribution and dipole orientation of molecules, implying site-dependent J_i in H_0. In addition, typical cavity couplings are very small ($g \sim 0.1J$) [44]. Figure 3(a) shows $T_{t'}$ as function of N for the same setup as in Fig. 2, with $\delta J = 0.2$ but for fixed $\Delta = g\sqrt{N} - J$. Without cavity ($g = 0$) T_{t_l} is exponentially suppressed, leading essentially to zero transmission for large N ($T_{t_s} < 10^{-6}$ for $N \gtrsim 400$ sites), as expected from Anderson-type localization [11]. However, the localized eigenstates of the system can be modified by the cavity [56–58]. Adding weak cavity couplings ($g = 0.05J, 0.1J, 0.2J$) already lifts the transport suppression and allows for a small but finite transmission even for systems with $N = 10^4$ sites. Consistent with the discussion above, in the collective strong coupling regime (right of the vertical lines) we find an universal algebraic behavior $T_{t'} \sim 1/N^2$. Interestingly, even in the weak coupling regime [left of vertical lines], i.e., when the two polariton peaks are not resolved, we find a small constant transmission orders of magnitudes above the cavity-free case.

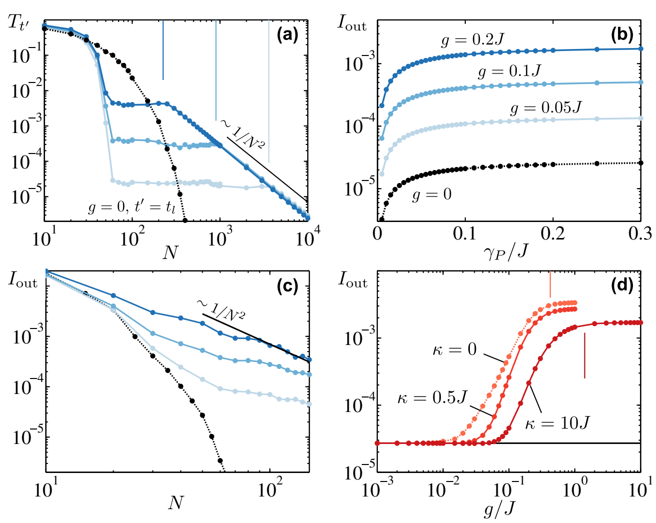

FIG. 3 (color online). Weakly coupled cavities. Blue lines and points denote $g/J = 0.05, 0.1, 0.2$ [light to dark]. $g = 0$ is shown as black dashed line with points. (a) T_{t_s} and T_{t_l} for wave-packet scattering [as in Fig. 2(a), $\Delta = g\sqrt{N} - J$] as function of N. We average over 200 disorder realizations in J_i ($\delta_J = 0.2J$). For $g = 0$ we plot T_{t_l}, for $g > 0$ T_{t_s}. T_{t_l} decreases exponentially with N. Cavity-assisted T_{t_s} becomes constant in the weak coupling regime and decays as $1/N^2$ (black solid line) in the collective strong coupling regime. Transitions between the regimes are indicated by vertical lines, marking $g\sqrt{N} = 3J$. (b) Steady state exciton current under incoherent pumping, $I_{\rm out}$, for $N = 50$ and in presence of disorder, spontaneous emissions, and dephasing (see text, $\gamma_{\rm out} = 2J$, $\gamma_P = 0.5J$). Small g already leads to massive enhancement. (c) Dissipation and disorder leads to $I_{\rm out} \propto \exp(-N)$ for $g = 0$, while for $g > 0$ $I_{\rm out}$ decays subexponentially. In the collective strong coupling regime, $I_{\rm out} \sim 1/N^2$ (black solid line). (d) $I_{\rm out}$ as a function of g. The crossover to the collective strong coupling regime shifts for $\kappa > 0$ ($N = 50$). Vertical lines: $g\sqrt{N} = 3J$ and $g\sqrt{N} = 10J$, respectively. (b)–(d): $\gamma_{\rm sp\ em} = 0.04J$, $\gamma_{\rm deph} = 0.9J$, and $\delta_J = 0.2J$, single disorder realization.

We note that in this Letter we deal with a one-dimensional situation, but we expect the main findings to also hold for dimensionalities $d > 1$. While relative improvement of T_{t_l} compared to the $g = 0$ case can decrease with increasing d (excitations can tunnel past impurities more easily), a finite T_{t_s} for $g = 0$ will also be impossible for $d = 2, 3$ because of the Lieb-Robinson bound [6]. In contrast, the cavity-photon mode occurs in any dimensions and thus one can expect the transmission mechanism to work in arbitrary dimensions. Since Anderson localization is also present in two dimensions, and in three dimensions below the mobility edge [59], an exponential improvement of transmission can be expected in such a situation.

Incoherent pumping setup.—We now consider the case (ii); i.e., we analyze steady-state currents $I_{\rm out}$ that develop under incoherent pumping of excitations ($\gamma_P, \gamma_{\rm out} > 0$). Spontaneous emission and dephasing are now included with $\gamma_{\rm sp\ em} = 0.04J$ and $\gamma_{\rm deph} = 0.9J$, respectively. Pump rates γ_P play the role similar to a "voltage" but for exciton currents. We plot $I_{\rm out}$-γ_P curves in Fig. 3(b). The figure shows that even small cavity couplings g can increase $I_{\rm out}$ by orders of magnitude compared to the cavity-free case. This finding is in stark contrast to previous works with exciton polaritons in multimode cavities [60] and constitutes one of the key results of this work.

Consistent with the wave-packet dynamics above, Fig. 3(c) shows that, for $g = 0$, $I_{\rm out}$ decreases exponentially with N, due to the various dissipative terms and the disorder. However, choosing $g = 0.05, 0.1, 0.2$, changes the currents dramatically: for $N = 150$ and $g = 0.2$ the collective strong coupling regime is barely reached $g\sqrt{N} \sim 2.5J$; nevertheless, remarkably, we find that $I_{\rm out}$, just as T_{t_s} above, already displays an algebraic $1/N^2$ decrease. The fact that the cavity enhancement of $I_{\rm out}$ is induced by a collective cavity coupling is further demonstrated in Fig. 3(d), where $I_{\rm out}$ is shown vs g, for a few values of κ. A sudden increase of $I_{\rm out}$ occurs when g exceeds a particular value (vertical lines). By inspection, we find that this indeed corresponds to the point where $g\sqrt{N}$ exceeds all other energy scales. Consistently, this point is shifted to larger values of g for large $\kappa = 10$.

Conclusion and outlook.—In this work, we have shown that both incoherent and coherent exciton transport in a spin chain can be dramatically enhanced by collective coupling to the structured vacuum field of a Fabry-Perot cavity. These results may be relevant for disordered organic semiconductors at room temperature, where exciton conduction may be ameliorated by orders of magnitude, as well as for artificial media made of Rydberg atoms, polar molecules, or cold ions at sub-mK temperatures. It is an exciting prospect to investigate whether strong coupling can also induce the ultrafast propagation of classical and quantum correlations [41,42,61]. While in 1D a modified density matrix renormalization group technique [38–40] might provide us with an answer, the higher dimensional situation could be a first example where only the artificially engineered quantum simulator setups can do so. Finally, a key open challenge not addressed here is to explore the physical mechanisms behind the enhancement of charge conductivity as reported in experiments [12].

We note that results for transport in disordered organic semiconductors related to those reported here have been independently obtained by Feist and Garcia-Vidal [62].

We thank T. W. Ebbesen, F. J. Garcia Vidal, J. Feist, L. M. Moreno, and H. Ritsch for fruitful discussion. J. S. acknowledges hospitality from the Institute for Theoretical Physics at the University of Innsbruck and the University of Strasbourg. This work was supported by the ERC-St Grant ColdSIM (Grant No. 307688), EOARD, and UdS via Labex NIE and IdEX, the JQI, FWF-ANR via BLUESHIELD, and RYSQ, the NSF PFC at the JQI and JILA (PIF-1211914 and PFC-1125844), as well as the Austrian Science Fund (FWF) via Project No. P24968-N27 (C. G). Computations utilized the HPC UdS, and the Janus supercomputer, supported by NSF (CNS-0821794), NCAR, and CU Boulder.

[1] P. Jurcevic, B. P. Lanyon, P. Hauke, C. Hempel, P. Zoller, R. Blatt, and C. F. Roos, Nature (London) **511**, 202 (2014).
[2] P. Richerme, Z.-X. Gong, A. Lee, C. Senko, J. Smith, M. Foss-Feig, S. Michalakis, A. V. Gorshkov, and C. Monroe, Nature (London) **511**, 198 (2014).
[3] M. Cheneau, P. Barmettler, D. Poletti, M. Endres, P. Schauß, T. Fukuhara, C. Gross, I. Bloch, C. Kollath, and S. Kuhr, Nature (London) **481**, 484 (2012).
[4] G. Günter, H. Schempp, M. Robert-de Saint-Vincent, V. Gavryusev, S. Helmrich, C. S. Hofmann, S. Whitlock, and M. Weidemüller, Science **342**, 954 (2013).
[5] S. Bose, Contemp. Phys. **48**, 13 (2007).
[6] E. H. Lieb and D. W. Robinson, Commun. Math. Phys. **28**, 251 (1972).
[7] B. Nachtergaele, Y. Ogata, and R. Sims, J. Stat. Phys. **124**, 1 (2006).
[8] S. R. Forrest, Nature (London) **428**, 911 (2004).
[9] G. D. Scholes and G. Rumbles, Nat. Mater. **5**, 683 (2006).
[10] S. M. Menke, W. A. Luhman, and R. J. Holmes, Nat. Mater. **12**, 152 (2013).
[11] P. W. Anderson, Phys. Rev. **109**, 1492 (1958).
[12] E. Orgiu, J. George, J. A. Hutchison, E. Devaux, J.-F. Dayen, B. Doudin, F. Stellacci, C. Genet, P. Samori, and T. W. Ebbesen, arXiv:1409.1900.
[13] D. G. Lidzey, D. D. C. Bradley, M. S. Skolnick, T. Virgili, S. Walker, and D. M. Whittaker, Nature (London) **395**, 53 (1998).
[14] D. G. Lidzey, D. D. C. Bradley, T. Virgili, A. Armitage, M. S. Skolnick, and S. Walker, Phys. Rev. Lett. **82**, 3316 (1999).
[15] R. J. Holmes and S. R. Forrest, Phys. Rev. Lett. **93**, 186404 (2004).
[16] D. M. Coles, P. Michetti, C. Clark, W. C. Tsoi, A. M. Adawi, J.-S. Kim, and D. G. Lidzey, Adv. Funct. Mater. **21**, 3691 (2011).
[17] T. Schwartz, J. A. Hutchison, C. Genet, and T. W. Ebbesen, Phys. Rev. Lett. **106**, 196405 (2011).
[18] S. Kéna-Cohen, S. A. Maier, and D. D. C. Bradley, Adv. Opt. Mater. **1**, 827 (2013).
[19] J. D. Plumhof, T. Stöferle, L. Mai, U. Scherf, and R. F. Mahrt, Nat. Mater. **13**, 247 (2014).
[20] R. Balili, V. Hartwell, D. Snoke, L. Pfeiffer, and K. West, Science **316**, 1007 (2007).
[21] E. Wertz, A. Amo, D. D. Solnyshkov, L. Ferrier, T. C. H. Liew, D. Sanvitto, P. Senellart, I. Sagnes, A. Lemaître, A. V. Kavokin *et al.*, Phys. Rev. Lett. **109**, 216404 (2012).
[22] M. Alloing, M. Beian, M. Lewenstein, D. Fuster, Y. González, L. Gonzlez, R. Combescot, M. Combescot, and F. Dubin, Europhys. Lett. **107**, 10012 (2014).
[23] H. Shirakawa, E. J. Louis, A. G. MacDiarmid, C. K. Chiang, and A. J. Heeger, J. Chem. Soc., Chem. Commun. 578 (1977).
[24] C. W. Tang and S. A. VanSlyke, Appl. Phys. Lett. **51**, 913 (1987).
[25] J. H. Burroughes, D. D. C. Bradley, A. R. Brown, R. N. Marks, K. Mackay, R. H. Friend, P. L. Burns, and A. B. Holmes, Nature (London) **347**, 539 (1990).
[26] C. Kim, P. E. Burrows, and S. R. Forrest, Science **288**, 831 (2000).
[27] A. Kahn, N. Koch, and W. Gao, J. Polym. Sci. **41**, 2529 (2003).
[28] H. Sirringhaus, M. Bird, and N. Zhao, Adv. Mater. **22**, 3893 (2010).
[29] A. C. Arias, J. D. MacKenzie, I. McCulloch, J. Rivnay, and A. Salleo, Chem. Rev. **110**, 3 (2010).
[30] J. A. Hutchison, A. Liscio, T. Schwartz, A. Canaguier-Durand, C. Genet, V. Palermo, P. Samorì, and T. W. Ebbesen, Adv. Mater. **25**, 2481 (2013).
[31] R. Löw, H. Weimer, J. Nipper, J. B. Balewski, B. Butscher, H. P. Büchler, and T. Pfau, J. Phys. B **45**, 113001 (2012).
[32] S. Bettelli, D. Maxwell, T. Fernholz, C. S. Adams, I. Lesanovsky, and C. Ates, Phys. Rev. A **88**, 043436 (2013).
[33] B. Yan, S. A. Moses, B. Gadway, J. P. Covey, K. R. A. Hazzard, A. M. Rey, D. S. Jin, and J. Ye, Nature (London) **501**, 521 (2013).
[34] K. R. A. Hazzard, B. Gadway, M. Foss-Feig, B. Yan, S. A. Moses, J. P. Covey, N. Y. Yao, M. D. Lukin, J. Ye, D. S. Jin *et al.*, Phys. Rev. Lett. **113**, 195302 (2014).

[35] D. Porras and J. I. Cirac, Phys. Rev. Lett. **92**, 207901 (2004).
[36] A. Friedenauer, H. Schmitz, J. T. Glueckert, D. Porras, and T. Schaetz, Nat. Phys. **4**, 757 (2008).
[37] J. M. Raimond, M. Brune, and S. Haroche, Rev. Mod. Phys. **73**, 565 (2001).
[38] G. Vidal, Phys. Rev. Lett. **93**, 040502 (2004).
[39] S. R. White and A. E. Feiguin, Phys. Rev. Lett. **93**, 076401 (2004).
[40] A. J. Daley, C. Kollath, U. Schollwöck, and G. Vidal, J. Stat. Mech. (2004) P04005.
[41] Z.-X. Gong, M. Foss-Feig, S. Michalakis, and A. V. Gorshkov, Phys. Rev. Lett. **113**, 030602 (2014).
[42] J. Schachenmayer, B. P. Lanyon, C. F. Roos, and A. J. Daley, Phys. Rev. X **3**, 031015 (2013).
[43] A. Rodríguez, V. A. Malyshev, G. Sierra, M. A. Martín-Delgado, J. Rodríguez-Laguna, and F. Domínguez-Adame, Phys. Rev. Lett. **90**, 027404 (2003).
[44] See Supplemental Material http://link.aps.org/supplemental/10.1103/PhysRevLett.114.196403 for details on scattering calculations as well as considerations and numbers for possible experiments, which includes Refs. [45–47].
[45] F. Robicheaux, J. V. Hernández, T. Topçu, and L. D. Noordam, Phys. Rev. A **70**, 042703 (2004).
[46] A. Stute, B. Casabone, P. Schindler, T. Monz, P. O. Schmidt, B. Brandstätter, T. E. Northup, and R. Blatt, Nature (London) **485**, 482 (2012).
[47] R. H. Lehmberg, Phys. Rev. A **2**, 883 (1970).
[48] E. I. Rashba, Synth. Met. **64**, 255 (1994).
[49] V. M. Agranovich and A. M. Kamchatnov, Chem. Phys. **245**, 175 (1999).
[50] D. Manzano, M. Tiersch, A. Asadian, and H. J. Briegel, Phys. Rev. E **86**, 061118 (2012).
[51] D. A. Ryndyk, R. Gutierrez, B. Song, and G. Cuniberti, arXiv:0805.0628 [Lect. Notes Phys. (to be published)].
[52] R. Landauer, IBM J. Res. Dev. **1**, 223 (1957).
[53] M. Büttiker, Y. Imry, R. Landauer, and S. Pinhas, Phys. Rev. B **31**, 6207 (1985).
[54] M. Biondi, S. Schmidt, G. Blatter, and H. E. Türeci, Phys. Rev. A **89**, 025801 (2014).
[55] P. Longo, P. Schmitteckert, and K. Busch, Phys. Rev. Lett. **104**, 023602 (2010).
[56] A. Biella, arXiv:1407.6562.
[57] A. Biella, F. Borgonovi, R. Kaiser, and G. L. Celardo, Europhys. Lett. **103**, 57009 (2013).
[58] G. Celardo, A. Biella, L. Kaplan, and F. Borgonovi, Fortschr. Phys. **61**, 250 (2013).
[59] E. Abrahams, P. W. Anderson, D. C. Licciardello, and T. V. Ramakrishnan, Phys. Rev. Lett. **42**, 673 (1979).
[60] V. M. Agranovich and Y. N. Gartstein, Phys. Rev. B **75**, 075302 (2007).
[61] M. Litinskaya, Phys. Rev. B **77**, 155325 (2008).
[62] J. Feist and F. J. Garcia-Vidal, preceding Letter, Phys. Rev. Lett. **114**, 196402 (2015).