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Nonequilibrium phononic first-order phase transition in a driven fermion chain

Mohsen Yarmohammadi,^{1,*} Marin Bukov,^{2,†} and Michael H. Kolodrubetz^{1,‡}

¹*Department of Physics, The University of Texas at Dallas, Richardson, Texas 75080, USA*

²*Max Planck Institute for the Physics of Complex Systems, Nöthnitzer Str. 38, 01187 Dresden, Germany*

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We study the direct laser drive of infrared-active phonons that are quadratically coupled to a spinless fermion chain. Feedback is incorporated by phonon dressing of the electronic dispersion, which enables effective non-linearities in the phonon dynamics. We uncover a first-order phase transition in the phononic steady state in which hysteretic effects allow either large or small phonon occupation depending on the drive protocol. We discuss the implications of these findings for probing phase transitions in real driven materials.

Introduction.—Despite a considerable ongoing effort to understand quantum systems, uncovering nonequilibrium phenomena without an apparent equilibrium analogue is one of the outstanding scientific challenges in modern condensed matter physics. Recent developments in the ultrafast dynamics of quantum materials have attracted interest in exploring novel nonequilibrium many-body phenomena [1–7]. Interacting light-matter systems with competing interactions emerged as a useful experimental platform since they allow to study the interplay between coupling to an environment and external drives [8–10]. Moreover, they play a central role in the study of photovoltaics [11, 12], light-induced phase transitions [13, 14], and laser processing [15].

Direct laser driving of an infrared (IR)-active phonon is the main light-matter interaction channel; it couples to the electron degree of freedom by modifying its environment. In a typical setup, phonons couple to the electron number operator in a nonlinear manner [16]. Over the last decade, a considerable effort has been devoted to studying the role of electron-phonon coupling on the dynamical properties of high-temperature cuprate superconductors [17, 18] and Mott or charge-density-wave insulators [19–22]. For a driven electron-phonon system to avoid heating (e.g., with the goal to store and transmit information), the latter needs to be counterbalanced by dissipation from undriven phonon modes. This has been studied in spin lattices [23–25], coupled quantum-electrodynamics cavities and circuits [26–28], lattice Rydberg atoms [29–31], driven-dissipative superfluids [32], nonlinear photonic modes [33, 34], etc. Although dissipation is widely captured in the dynamics of local observables, a detailed understanding of the process is still an outstanding challenge.

Studies of interacting systems, exposed to an external laser field and coupled to a thermal bath, have until recently mostly focused on the “bare” dispersion of degrees of freedom in dissipation processes. However, it remains to be understood how the backaction of transient excitations can be incorporated into the dissipation process. To the best of our knowledge, this has not been addressed so far in phononically driven materials. This mainly requires a deep scan of the “dressed” electronic dispersion

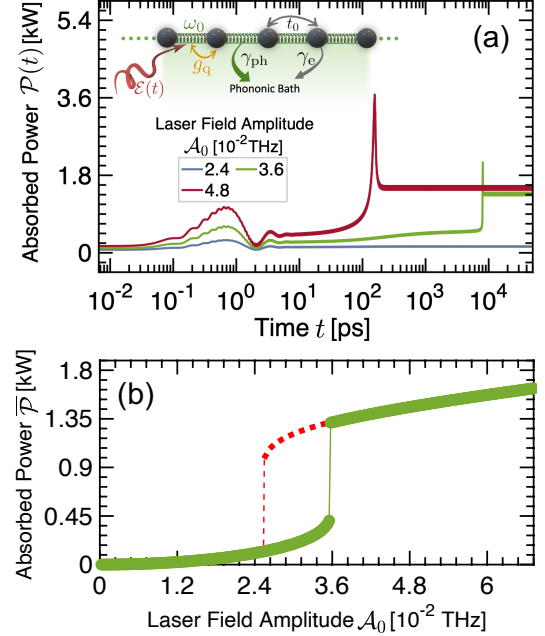


FIG. 1. Nonequilibrium first-order phase transition in a phononically driven fermion chain. (a) Drive-period-averaged time evolution of absorbed power for a driven-dissipative chain of spinless fermions, sketched in the inset, starting from the ground state at half filling. The system reaches a nonequilibrium steady state (NESS) which undergoes a first-order phase transition at laser amplitude $A_0 \approx 3.6 \times 10^{-2}$ THz [solid line in (b)] due to phononic modification of the electronic dispersion. The dashed line in (b) shows a second stable NESS which can be created by modifying either the initial state or drive protocol. The parameters are chosen to be $\omega = 4.44$ THz, $\omega_0 = 4.8$ THz, $g_q = 9.6$ THz, $\gamma_{ph} = 0.24$ THz, and $\gamma_e = 0.0024$ THz; see text for details.

through a time-dependent damping rate, which results in a novel physical insight.

In this Letter, we show that a driven-dissipative fermion chain with a dressed dispersion, exhibits a *dynamical first-order phase transition* due to local depopulation of the electrons in the spirit of quadratic electron-phonon coupling (QEPC). One experimental signature of the transition is a sharp spike in the long time-evolution of the absorbed power, accompanied by a discontinu-

ity in its value in the long-time nonequilibrium steady state (NESS), as shown in Fig. 1. We analyze and identify accessible parameter regimes to provide a clear path for experimental verification.

Model.—Consider an infinite half-filled chain of spinless fermions with periodic boundary conditions, as illustrated in Fig. 1(a), inset. We drive the chain by a continuous field coupled to the IR-active optical phonons, which illuminates the entire system until it reaches a NESS. The model Hamiltonian reads [35–37]

$$\mathcal{H}(t) = -t_0 \sum_{\ell} (c_{\ell}^{\dagger} c_{\ell+1} + \text{H.c.}) + \omega_0 \sum_{\ell} a_{\ell}^{\dagger} a_{\ell} + g_q \sum_{\ell} (a_{\ell}^{\dagger} + a_{\ell})^2 (c_{\ell}^{\dagger} c_{\ell} - 1/2) + \mathcal{E}(t) \sum_{\ell} (a_{\ell}^{\dagger} + a_{\ell}), \quad (1)$$

where $c_{\ell}^{\dagger}$ (c_{ℓ}) and $a_{\ell}^{\dagger}$ (a_{ℓ}) are respectively the electron and phonon creation (annihilation) operators at lattice site ℓ ; t_0 is a constant hopping amplitude setting the energy scale of our model, ω_0 is the optical phonon frequency, $c_{\ell}^{\dagger} c_{\ell} - 1/2$ is the electron number operator relative to half filling, and g_q is the strength of QEPC. In the laser-phonon coupling, the laser field is described by $\mathcal{E}(t) = \mathcal{A}_0 \cos(\omega t)$ with frequency ω and amplitude $\mathcal{A}_0$. In momentum space, we have the bare electronic dispersion $\omega_k = -2t_0 \cos k$ and a k -independent phonon frequency; the full k -space Hamiltonian and details of the model are provided in Sec. S1 of the Supplemental Material [38].

We focus our simulations on the dynamics of the approximate dispersionless phonon since, in most materials, the average ω_0 is much larger than the phonon bandwidth. Due to the relatively long wavelength of the drive phonons compared to the lattice spacing, the phonon response is dominated by the zero-momentum mode, a_0 . We also neglect linear and higher electron-phonon coupling effects, since in target centrosymmetric structures [36, 37] the dominant coupling is quadratic; the contribution of linear electron-phonon coupling is negligible due to a vanishing average of linear oscillations. For comparison to experiments on materials with nonlinear electron-phonon coupling, we choose to consider a representative hopping energy $t_0 = 10 \text{ meV} \approx 2.4 \text{ THz}$, which can also be chosen to be similar to (Pb,Bi)₂Sr₂CaCu₂O₈, YBa₂Cu₃O_{6+x}, and K₃C₆₀ [36, 39–41] without qualitatively modifying the results.

The system is coupled to a phononic bath (given by independent phonon modes not subject to the drive) which allows the formation of a NESS. To model the dynamics of the dissipation process, we use the Lindblad master equation for an arbitrary observable $O(t)$:

$$\begin{aligned} \langle \dot{O} \rangle(t) &= i \langle [\mathcal{H}, O(t)] \rangle \\ &+ \frac{1}{2} \sum_{\ell} \gamma_{\ell} \left\langle [\mathcal{L}_{\ell}^{\dagger}, O(t)] \mathcal{L}_{\ell} + \mathcal{L}_{\ell}^{\dagger} [O(t), \mathcal{L}_{\ell}] \right\rangle, \end{aligned} \quad (2)$$

where the summation ℓ describes a bounded set of operators for a separable (system-bath) Hilbert space.

The relevant observables are net electron density $n_e(t) = L^{-1} \sum_k \langle c_k^{\dagger} c_k \rangle(t)$, phonon displacement $q_{\text{ph}}(t) = L^{-1/2} \langle a_0^{\dagger} + a_0 \rangle(t)$, momentum $p_{\text{ph}}(t) = L^{-1/2} \langle i(a_0^{\dagger} - a_0) \rangle(t)$, occupation $n_{\text{ph}}(t) = L^{-1} \langle a_0^{\dagger} a_0 \rangle(t)$, as well as squeezing-related observables $\mathcal{Q}_{\text{ph}}(t) = L^{-1} \langle (a_0^{\dagger} a_0^{\dagger} + a_0 a_0) \rangle(t)$ and $\mathcal{P}_{\text{ph}}(t) = L^{-1} \langle i(a_0^{\dagger} a_0^{\dagger} - a_0 a_0) \rangle(t)$. We consider conventional phenomenological relaxation processes given by conventional local jump operators $\mathcal{L}_{\ell} = a_0^{\dagger}, a_0, c_k^{\dagger}$, and c_k with rates $\gamma_{\text{ph}} \mathcal{N}_0$, $\gamma_{\text{ph}}[1 + \mathcal{N}_0]$, $\gamma_e \tilde{\mathcal{N}}_k$, and $\gamma_e(1 - \tilde{\mathcal{N}}_k)$, respectively [25, 42–44], where $\mathcal{N}_0 = 1/[e^{\omega_0/k_B T} - 1]$ is the average number of phonons in equilibrium and $\tilde{\mathcal{N}}_k = 1/[e^{\tilde{\omega}_k/k_B T} + 1]$ is the equilibrium electron filling at momentum k .

Crucially, in defining the equilibrium electron occupation, we use the *dressed* electron dispersion $\tilde{\omega}_k(t) = \omega_k + g_q q_{\text{ph}}^2(t)$ rather than the bare dispersion. Therefore, as the phonon fluctuations q_{ph}^2 become large, they shift the chemical potential of the electrons downward, modifying the NESS electron density in the region where the laser drive couples. As we will see, this dressed dispersion gives rise to a nonlinearity that is crucial in understanding the steady state. Note that, by allowing relaxation of the electron density due to a dressed thermal bath, we are assuming that the electrons are connected to some external reservoir, such as metallic leads or the undriven region of the sample. Moreover, we assume that the chain is in touch with a cooling apparatus [25] to remediate the heating created by the continuous drive. Although our results are robust to the parameters chosen, we consider the experimentally relevant $\omega_0 = 4.8 \text{ THz}$, $\gamma_{\text{ph}} = 0.24 \text{ THz}$, and $\gamma_e = 0.0024 \text{ THz}$ throughout the text, unless otherwise specified. While certain parameters are rationally related, we have confirmed that resonances are irrelevant to the phenomena we describe below.

We initialize the system in its ground state (electrons at half filling, phonons in the vacuum state), and let it evolve in time following the Lindblad formalism [42, 43]. The expectation value of the phonon occupation is proportional to the number of sites L , thus, the relative quantum fluctuations which are proportional to $1/\sqrt{L}$ tend to zero in the thermodynamic limit ($L \rightarrow \infty$). Additionally, the phonon only couples to electrons via the site-averaged electron density, averaging out on-site electronic fluctuations. We, therefore, employ a mean-field approximation [25] to decouple the QEPC acting on the phonon and electron, factorizing expectation values according to $\langle A_{\text{ph}} B_e \rangle \approx \langle A_{\text{ph}} \rangle \langle B_e \rangle$. The equations then reduce to

$$\dot{q}_{\text{ph}}(t) = +\omega_0 p_{\text{ph}}(t) - \frac{\gamma_{\text{ph}}}{2} q_{\text{ph}}(t), \quad (3a)$$

$$\dot{p}_{\text{ph}}(t) = -\frac{\tilde{\omega}_0^2(t)}{\omega_0} q_{\text{ph}}(t) - 2\mathcal{E}(t) - \frac{\gamma_{\text{ph}}}{2} p_{\text{ph}}(t), \quad (3b)$$

$$\begin{aligned} \dot{n}_{\text{ph}}(t) &= -\mathcal{E}(t) p_{\text{ph}}(t) - 2g_q [n_e(t) - 1/2] \mathcal{P}_{\text{ph}}(t) \\ &\quad - \gamma_{\text{ph}} n_{\text{ph}}(t), \end{aligned} \quad (3c)$$

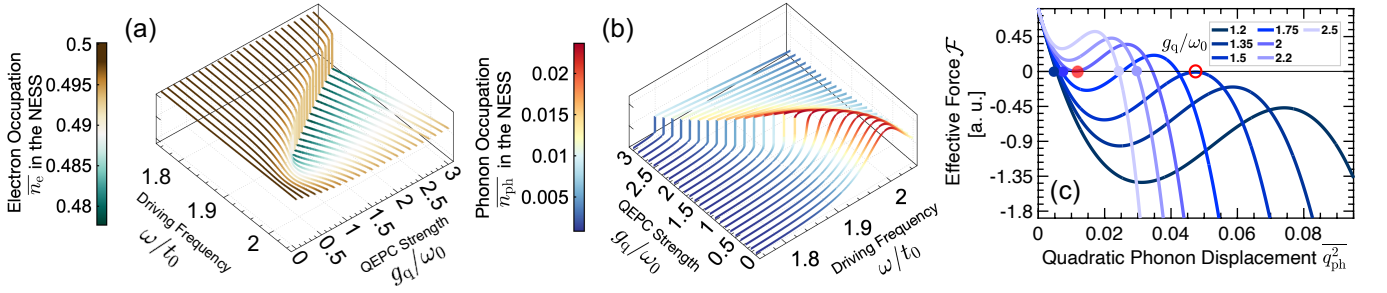


FIG. 2. **Origin of nonequilibrium first-order phase transition.** Dynamical response of dressed (a) electron and (b) phonon occupations in the nonequilibrium steady state for various quadratic electron-phonon couplings and driving frequencies (axes are swapped in (b) to improve visibility of jumps at the phase transition). Depopulation of the electrons leads to shifting of the phonon resonance, a nonlinearity which enables a first-order phase transition in the NESS at strong g_q . (c) Nonequilibrium effective force (see text) for various g_q at driving frequency $\omega = 4.44$ THz, which is below the bare phonon frequency. Shaded blue dots indicate the numerical steady state. Red dots indicate the two predicted phase transitions (see text), with bistability for $3/2 < g_q/\omega_0 < 2$. Parameters are $\omega_0 = 4.8$ THz, $\mathcal{A}_0 = 0.036$ THz, $\gamma_{ph} = 0.24$ THz, and $\gamma_e = 0.0024$ THz.

$$\begin{aligned} \dot{Q}_{ph}(t) = & + \left[\omega_0 + \frac{\tilde{\omega}_0^2(t)}{\omega_0} \right] \mathcal{P}_{ph}(t) + 2\mathcal{E}(t) p_{ph}(t) \\ & - \gamma_{ph} Q_{ph}(t), \end{aligned} \quad (3d)$$

$$\begin{aligned} \dot{P}_{ph}(t) = & - \left[\omega_0 + \frac{\tilde{\omega}_0^2(t)}{\omega_0} \right] Q_{ph}(t) - 2\mathcal{E}(t) q_{ph}(t) \\ & - 4g_q [2n_{ph}(t) + 1/L] [n_e(t) - 1/2] - \gamma_{ph} \mathcal{P}_{ph}(t), \end{aligned} \quad (3e)$$

$$\dot{n}_{e,k}(t) = -\gamma_e [n_{e,k}(t) - \tilde{N}_k(t)], \quad (3f)$$

where the electron occupation shifts the effective phonon frequency to $\tilde{\omega}_0(t) = \omega_0 \sqrt{1 + 4g_q[n_{e,k}(t) - 1/2]/\omega_0}$ [45].

We consider $T = 0$, such that $\tilde{N}_k$ equals 1 (filled states) for $|k| < \tilde{k}_F$ and 0 (empty) otherwise, where $\tilde{k}_F$ is the Fermi momentum for the dressed dispersion $\tilde{\omega}_k$. Despite the complexity of these equations, they can be solved efficiently on a computer; we do so for chain lengths up to $L = 1001$ sites.

As the QEPC g_q plays an essential role in determining the dynamics of the model, it is worth noting that the renormalization of the phonon frequency in our model, i.e., $\tilde{\omega}_0 = \omega_0 \sqrt{1 + 4g_q[\langle n_{e,\ell} \rangle - 1/2]/\omega_0}$ does not lead to instabilities in the simulation, in contrast to other works [36, 37] where the system is unstable for $|g_q| > \omega_0/2$. The reason for this is our mean-field-type approximation, which replaces the quantized on-site electron occupation by an average value that only weakly deviates from $1/2$. Crucially, this lack of instability is also physical, suggesting that our results will survive fluctuations. This is because the constraint $|g_q| > \omega_0/2$ comes from the assumption of dispersionless Einstein phonons; in reality, there will always be some finite dispersion, causing excitations of the phonons to be spread over at least a few sites. If the minimal spatial dispersion is ℓ sites, the instability threshold should be roughly increased by a factor of ℓ , which brings it well above the value for g_q in real materials. Hence, our approximation allows us to simulate the model for arbitrary QEPC strength.

Results.—In addition to microscopic quantities such as

electron and phonon occupation, we consider the NESS energy flow from the drive into the final stage of dissipation among various degrees of freedom. Of particular importance is the absorbed power

$$\mathcal{P}(t) = -w a \rho \omega_0 \mathcal{E}(t) p_{ph}(t). \quad (4)$$

We use parameters from the YBCO sample [46], namely thickness $w = 10$ nm, area $a = 1$ mm², and molar density $\rho \approx 0.007$ mol.cm⁻³. This power is related to experimentally measurable quantities such as reflectance [47–54].

Upon quenching on the drive and solving the dynamics, the system evolves into a NESS as seen in Fig. 1(a). For the majority of parameters, it takes around 600 ps to reach the NESS, which is within the current technological capabilities of laser sources in ultrafast experiments. In the presence of a strong QEPC, as laser amplitude $\mathcal{A}_0$ is increased, the smooth transient process of the time-evolved absorbed power suddenly becomes a sharp peak at a critical value of $\mathcal{A}_0 = 0.036$ THz. At this point, observables including the absorbed power undergo a non-analytic change and a *first-order phase transition* occurs, which comes from the depopulation of the electronic modes via feedback from the phonons in the dissipation process, i.e., from $\tilde{N}_k(t)$ in Eq. (3f). This can also be confirmed by focusing on the NESS in Fig. 1(b), i.e., by averaging an observable $O(t)$ over one drive period, denoted by $\bar{O}$. As we show later, modifying either the initial state or the drive protocol leads to a different NESS, as shown by the dashed red line in Fig. 1(b) and supported by the data in Fig. 4.

To explain the origin of phase transition and to check its robustness against changing model parameters, we continue by discussing the response of the electrons and phonons to changing drive frequency ω and QEPC g_q [Figs. 2(a) and 2(b)]. In the absence of QEPC, there is a resonant peak at $\omega = \omega_0$, more clearly visible in the phonon occupation. As g_q is increased, the peak shifts to a lower frequency due to the depopulation of the electron density via feedback from the phonons. Surprisingly,

above a critical value of g_q , the smooth peak suddenly becomes a sharp jump, suggestive of a first-order phase transition.

The main origin of the phase transition is dissipation-induced nonlinearity. The dressed electron dispersion $\tilde{\omega}_k(t) = \omega_k + g_q q_{\text{ph}}^2(t)$ will yield an average shift of the electron chemical potential given by $\overline{g_q q_{\text{ph}}^2}$ (note that $\overline{q_{\text{ph}}} = 0$). Linearizing the electron density around the Fermi surface via $L^{-1} \sum_k \tilde{N}_k(t) = \tilde{k}_F(t)/\pi$, where $\tilde{k}_F(t) = \pi/2 - g_q q_{\text{ph}}^2(t)/2t_0$, we predict a steady state electron density of

$$\overline{n_e} = \frac{1}{2} - \frac{g_q \overline{q_{\text{ph}}^2}}{2\pi t_0}. \quad (5)$$

Since the steady state to which the electrons attempt to relax to depends on the phonon state, whose dynamics, in turn, depend on the electron density, this results in effective nonlinearities in the dynamics. This backaction was not considered in previous works [36, 37] because the electron was assumed to relax to its undriven ground state. While reasonable for linear electron-phonon coupling, the presence of a finite $\overline{q_{\text{ph}}^2}$ makes this backaction crucial for QEPC [55].

Correctly accounting for this backaction is our major contribution to the model which, as we have seen, produces significant effects on the dynamics. It is also important to note that $q_{\text{ph}}^2(t)$ can be experimentally measured through the intensity of a diffraction peak in femtosecond time-resolved X-rays [56], enabling another path to measure the effects of this backaction.

To microscopically interpret the observed phase transition, we consider the system near its steady-state; we assume that the phonons synchronize with the drive, with the dominant Fourier component of $q_{\text{ph}}(t) = q_1 e^{i\omega t} + \text{c.c.}$ and similarly for p_{ph} . Away from the NESS, q_1 slowly evolves towards a stationary point, which can be obtained by setting $\dot{q}_1 = \dot{p}_1 = 0$. Solving Eqs. (3a) and (3b) for this Fourier mode, we obtain

$$i\omega q_1 = \omega_0 p_1 - \frac{\gamma_{\text{ph}}}{2} q_1, \quad (6a)$$

$$i\omega p_1 = -\omega_0 q_1 - 4q_q[\overline{n_e} - 1/2]q_1 - 2\mathcal{A}_0 - \frac{\gamma_{\text{ph}}}{2} q_1. \quad (6b)$$

Using Eq. (5) and the relation $\overline{q_{\text{ph}}^2} = |q_1|^2/2$, this gives a cubic equation

$$\mathcal{F} = -\frac{4g_q^4\omega_0^2}{\pi^2 t_0^2} \overline{q_{\text{ph}}^2}^3 - \frac{4g_q^2\omega_0}{\pi t_0} \left(\omega^2 - \omega_0^2 + \frac{\gamma_{\text{ph}}^2}{4} \right) \overline{q_{\text{ph}}^2}^2 - \left(\gamma_{\text{ph}}^2 \omega^2 + (\omega^2 - \omega_0^2)^2 \right) \overline{q_{\text{ph}}^2} + 2\mathcal{A}_0^2 \omega_0^2, \quad (7)$$

with $\mathcal{F} = 0$ in the NESS. Perturbing away from the $\mathcal{F} = 0$ solution, this cubic equation may be thought of as a generalized force; stable attractors have $d\mathcal{F}/d\overline{q_{\text{ph}}^2} < 0$.

As seen in Fig. 2(c), the force curve shifts depending on various parameters and appears to have two separate phase transitions (marked by red circles) from having a unique NESS to a regime with two distinct stable NESSs. The first phase transition (transparent circle) is visible in

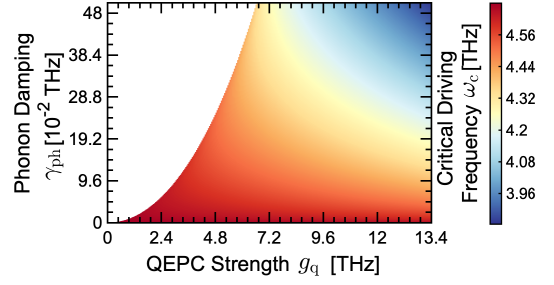


FIG. 3. Dynamical damping-coupling constraint for the phase transition. Critical driving frequency as a function of phonon damping rate and strength of quadratic electron-phonon coupling at $\omega_0 = 4.8$ THz, $\mathcal{A}_0 = 0.036$ THz, and $\gamma_e = 0.0024$ THz. The transition occurs down to arbitrarily small values of g_q and γ_{ph} . The separatrix, below which dynamical first-order phase transitions occur, is parabolic in the $\gamma_{\text{ph}} - g_q$ plane.

our numerical data since, starting from a vacuum state, a jump from small to large $\overline{q_{\text{ph}}^2}$ matches our expectations. The second phase transition (open circle) is not seen in the data, but could be realized via protocols similar to those in Fig. 4. This bifurcation transition with a cubic equation for the effective force is reminiscent of the magnetization in a first-order Ising phase transition. However, we emphasize that our phase transition occurs not in equilibrium, but rather out of equilibrium in the driven NESS.

One may naively conclude from Fig. 2 that the phase transition requires extremely large QEPCs, which can be challenging to find experimentally. However, by solving the cubic equation, we find the following analytical expression for the critical driving frequency:

$$\omega_c = \sqrt{\omega_0^2 - \frac{2\sqrt{3}\gamma_{\text{ph}}\omega_0}{3}} + \frac{W}{2}, \quad (8a)$$

$$W = \sqrt{\gamma_{\text{ph}}^4 - \frac{\pi t_0 \tilde{b}}{3g_q^2\omega_0} - 4\gamma_{\text{ph}}^2\omega_0^2 - \frac{96g_q^4\omega_0^4\mathcal{A}_0^2}{\pi^2 t_0^2 \tilde{b}}}, \quad (8b)$$

where $\tilde{b} = g_q^2\omega_0\gamma_{\text{ph}}(3\gamma_{\text{ph}} - 4\sqrt{3}\omega_0)/\pi t_0$. Importantly, this shows that a phase transition persists down to arbitrarily small g_q , as seen in the parametric plot of γ_{ph} vs. g_q in Fig. 3. Intriguingly, even with 1% damping of phonon energy ($\gamma_{\text{ph}} \approx 0.048$ THz) to the phononic bath, governed by the Lindemann criterion [57], the system still features the phase transition for moderate QEPC $g_q = \omega_0/2 = t_0$. The region in which the transition can be observed increases with the phonon frequency.

The generalized force in Fig. 2(c) suggests that two stable equilibria exist at large g_q or small ω , yet only one NESS is seen in our simulation for each value of the parameters. This is the attractor of the equations of motion starting from our initial ground state. To realize the other NESS, a different initial state must be prepared. Taking the cue from hysteresis near equilibrium first-order phase transitions, we suggest that this

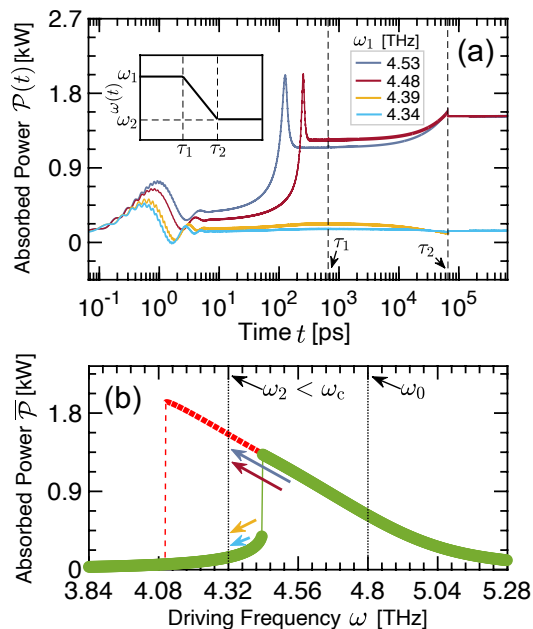


FIG. 4. **Chirp protocol for realizing the alternative steady state.** (a) Absorbed power in presence of a linearly chirped electric field $\mathcal{E}(t) = \mathcal{A}_0 \cos(\omega(t)t)$ with $\omega(t) = \omega_1 + [(\omega_2 - \omega_1)(t - \tau_1)/(\tau_2 - \tau_1)]$ (inset) for fixed $\omega_2 = 4.32$ THz $< \omega_c = 4.44$ THz. (b) Depending on whether ω_1 is less than or greater than ω_c , the system ends up in a different branch of the steady state. Black arrows point to the positions of the frequencies ω_0 and ω_2 (black dotted vertical lines); color-coded arrows indicate the value of ω_1 , cf. legend in (a). Parameters are $\omega_0 = 4.8$ THz, $\mathcal{A}_0 = 0.036$ THz, $\gamma_{ph} = 0.24$ THz, and $\gamma_e = 0.0024$ THz.

may be done via slowly ramping one of the parameters, such as the frequency (a “chirp” protocol). We consider a linearly chirped electric field $\mathcal{E}(t) = \mathcal{A}_0 \cos(\omega(t)t)$ with $\omega(t) = \omega_1 + (\omega_2 - \omega_1)(t - \tau_1)/(\tau_2 - \tau_1)$, as shown in Fig. 4(a), inset. Such frequency ramps are important throughout ultrafast spectroscopy [58, 59], and can thus be implemented experimentally. Starting with frequency ω_1 at time τ_1 sufficiently large to reach the NESS, we slowly ramp to the final value $\omega_2 < \omega_c$ at time τ_2 . If $\omega_1 < \omega_c$ as well, no phase transition is crossed by this ramp and the system simply reaches the original NESS. However, if $\omega_1 > \omega_c$, the system instead stays in the upper NESS with large q_{ph}^2 , as shown in Fig. 4(b). Similar hysteretic preparation of the alternative NESS can be accomplished via slowly ramping other experimental parameters, such as the drive amplitude $\mathcal{A}_0$ in Fig. 1(b).

Conclusion/Outlook.—We have uncovered a novel nonequilibrium phase transition of a driven-dissipative fermion chain coupled to a phonon mode. By dynamically tuning the dissipation processes through the backaction of phononic excitations on the electronic dispersion, we find a robust phase transition from having a unique nonequilibrium steady state to having two stable steady states. Since the effect is favored by a quasi-equilibrium

electron density (less than 2% deviation from a half-filled chain in equilibrium), it should be prevalent in any fermion chain. Our findings motivate future nonequilibrium spectroscopy experiments to seek photo-induced phononic phase transitions in driven quantum materials such as driven superconductors, where the modified phononic steady state may enable a nonequilibrium pathway to controlling superconductivity [60, 61]. Since such strong phononic steady state tends to be inhibited by the feedback mechanism that we consider, nonequilibrium protocols such as the chirp may be required to push these systems into the desired NESS. This phase transition will likely also be possible in the presence of an electromagnetic cavity, which has been used in recent work to control heating in driven materials [62–64].

Finally, we note that these nonequilibrium phase transitions not restricted to 1D, and are in fact more likely in higher dimensions, where mean-field theory and expansions around the Fermi sea are even more well-justified. Indeed, the phenomenon we reported primarily comes from a mean-field treatment of the average density, whose weakly correlated behavior is insensitive to dimensionality up to an overall prefactor in the density of states.

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* mohsen.yarmohammadi@utdallas.edu

† mgbukov@pks.mpg.de

‡ mkolodru@utdallas.edu

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