PHYSICAL REVIEW B **94**, 041104(R) (2016)

Parity-time symmetry-breaking mechanism of dynamic Mott transitions in dissipative systems

Vikram Tripathi,[1,2] Alexey Galda,[1] Himadri Barman,[2] and Valerii M. Vinokur[1]

[1]*Materials Science Division, Argonne National Laboratory, 9700 S. Cass Avenue, Argonne, Illinois 60439, USA*

[2]*Department of Theoretical Physics, Tata Institute of Fundamental Research, Homi Bhabha Road, Mumbai 400005, India*

(Received 4 November 2015; revised manuscript received 13 June 2016; published 5 July 2016)

We describe the critical behavior of the electric field-driven (dynamic) Mott insulator-to-metal transitions in dissipative Fermi and Bose systems in terms of non-Hermitian Hamiltonians invariant under simultaneous parity ($\mathcal{P}$) and time-reversal ($\mathcal{T}$) operations. The dynamic Mott transition is identified as a $\mathcal{PT}$ symmetry-breaking phase transition, with the Mott insulating state corresponding to the regime of unbroken $\mathcal{PT}$ symmetry with a real energy spectrum. We establish that the imaginary part of the Hamiltonian arises from the combined effects of the driving field and inherent dissipation. We derive the renormalization and collapse of the Mott gap at the dielectric breakdown and describe the resulting critical behavior of transport characteristics. The obtained critical exponent is in an excellent agreement with experimental findings.

DOI: 10.1103/PhysRevB.94.041104

Non-Hermitian $\mathcal{PT}$-symmetric quantum Hamiltonian models introduced in the seminal work of Bender and Boettcher [1] offer a foundation for the description of nonequilibrium steady states of dissipative quantum systems [2–5]. The basic property of non-Hermitian $\mathcal{PT}$-symmetric models is that their eigenstates exhibit a continuous $\mathcal{PT}$ symmetry-breaking phase transition when the strength of the external nonconservative driving force exceeds a certain threshold value. Below this threshold, i.e., in the regime of *unbroken* $\mathcal{PT}$ symmetry, the energy eigenvalues are real, while above it the energy spectrum acquires an imaginary part.

The $\mathcal{PT}$-symmetric models arise across the entire nonequilibrium physics and describe optical waveguides [6], electric RLC circuits [7], microwave cavities [8], and superconducting wires [9], to name a few. Here we focus on a theory of the electric field- or current-driven Mott metal-insulator transitions (MIT) as a $\mathcal{PT}$ symmetry-breaking phenomenon. The interest in the nonequilibrium MIT is motivated by both the intellectual appeal of understanding dynamic instabilities in quantum many-body strongly correlated systems and the high technological promise of Mott systems as a platform for switching devices in emergent electronics [10]. There have been tantalizing reports of field-driven Mott MIT in VO_2 [11,12], $La_{2-x}Sr_xNiO_4$ [13], one-dimensional Mott insulators $Sr_2CuO_3/SrCuO_2$ [14], and organic compounds [15], yet the critical behavior at the dynamic Mott transition remained unexplored. In a recent experimental breakthrough [16], the current-driven Mott transition has been observed in a system of vortices pinned by a periodic array of proximity coupled superconducting islands. Notably, the revealed critical behavior of the dynamic resistance near the dynamic Mott critical point appeared to belong to the liquid-gas transition universality class. Here we propose the $\mathcal{PT}$ symmetry-breaking mechanism of the dynamic MIT and derive the critical behavior that perfectly agrees with the experimental findings.

There has been a remarkable progress in unearthing the mechanism of the field-driven breakdown of the Mott insulator, which was identified as the Landau-Zener-Schwinger (LZS) process of generation of free particle-hole excitations by an external driving field [17–19]. The remaining puzzle concerns the description of the collapse of the Mott gap at the transition. As we show below, this can be achieved by taking into account dissipation processes. A recent numerical study [20] that included dissipation still did not address the critical behavior. An important step towards including the dissipation effects into the picture was taken in [21] via the Bethe ansatz treatment of a half-filled Hubbard chain subject to a constant imaginary gauge field, the approach resembling the delocalization transition induced in a system of noninteracting vortices by an imaginary vector potential [22]. In an intriguing parallel development in high-energy physics, a numerical treatment of the Schwinger mechanism in scalar electrodynamics revealed mass renormalization due to thermalization of produced particles [23]. Dissipation-assisted enhancement of LZS tunneling is also known to occur in noninteracting systems [24].

Here we address the challenge of description of the collapse of the Mott gap at the transition. We develop a theory of the electric field-driven MIT based on the concept of the $\mathcal{PT}$ symmetry breaking. We show that it is the applied electric field which, in the presence of dissipation, generates an imaginary part of the system's Hamiltonian while retaining its $\mathcal{PT}$ symmetry. We consider fermionic and bosonic systems that undergo the transition and identify their MITs as $\mathcal{PT}$ symmetry-breaking phase transitions. For a half-filled Hubbard chain we adopt the Bethe ansatz approach [21] and obtain the critical scaling of the Mott gap Δ with driving field $F, \Delta \sim (F_c - F)^{1/2}$. Then we find the probability $P \sim e^{-2\gamma}$ for the LZS dielectric breakdown with the LZS tunneling parameter $\gamma \sim (F_c - F)^{3/2}$. For a two-dimensional model we employ a dynamical mean field theory (DMFT) approach with an iterative perturbation theory (IPT)-based impurity solver [25] and find a critical scaling $\Delta \sim (F_c - F)^{0.78\pm0.03}$. For the vortex (bosonic) Mott transitions driven by the current I, we derive $\Delta \sim (I_c - I)^{1/2}$ scaling of the spectral gap, and $\gamma \sim (I_c - I)^{3/2}$ collapse of the LZS parameter, in an excellent accord with experimental results of Ref. [16].

The model. Let $|0\rangle$ be the ground state of an interacting quantum system, $|1\rangle$ be the lowest excited state, and $\Delta = E_1 - E_0$ be the spectral gap, with E_0 and E_1 being the eigenvalues for the ground and lowest excited state, respectively. Within the LZS framework, the electric field-induced probability for the $|0\rangle \to |1\rangle$ transition is given by the Landau-Dykhne formula [26], $P = |\langle 0|1\rangle|^2 \sim e^{-2\gamma}$, where

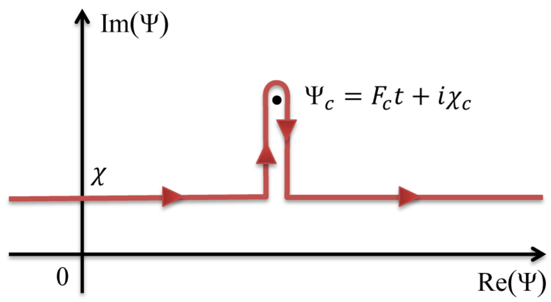

FIG. 1. Integration contour in the complex Ψ plane for calculating the Landau-Zener-Schwinger transition probability. The point $\Psi_c = F_c t + i\chi_c$ is a degeneracy point where the energy gap, $\Delta \equiv E_1 - E_0$, closes. Only the parts of the contour parallel to the imaginary axis contribute to the transition probability. For finite dissipation, the vertical part of the contour begins from a nonzero value of χ, whose implicit dependence on the dissipation and driving field is obtained in the main text.

$\gamma = (1/\hbar)\mathrm{Im}\int_{-\infty}^{\infty} dt\, \Delta[\Psi(t)]$, and Ψ is the time-dependent phase factor related to the driving field, F. For electrons hopping along a 1D lattice under the constant electric field, we choose the gauge where the driving field is the time derivative of the vector potential. We replace the integration over time by that over complex $\Psi = Ft + i\chi(F)$, where the imaginary component χ, shown below to be directly responsible for energy gap renormalization, originates from the effects of dissipation and is an odd function of F. Assuming that $\Psi(F)$ is differentiable and that at some critical $\chi_c \equiv \chi(F_c)$ the gap completely closes, we find that the imaginary part of the integral over Ψ comes from the degeneracy point $i\chi_c$, see Fig. 1:

$$\gamma \sim \frac{1}{F}\mathrm{Re}\int_{\chi}^{\chi_c} d\chi'[E_1(\chi') - E_0(\chi')]. \quad (1)$$

In nondissipative models, the Landau-Dykhne formula (1) reduces to the Landau-Zener result, $\gamma \sim \Delta^2/vF$, where $v = |d\Delta/dt|/F$ is the speed of convergence of E_1 and E_2 upon variation of Ψ, which is assumed to be constant (i.e., independent of Δ) in the usual Landau-Zener analysis. Writing $\gamma = F_{\mathrm{th}}/F$, we identify the threshold field for the LZS tunneling, $F_{\mathrm{th}} = \Delta^2/v$. The same relation has been obtained in Ref. [17] for the half-filled Hubbard chain.

Before turning to our major theme of how the external drive and dissipation induce the imaginary component of the vector potential, note that Eq. (1) implies that the Landau-Zener factor is determined completely by the imaginary part of Ψ. Hence in the following analysis we will be setting the real part of Ψ be zero. To see that such a choice does not lead to any loss of generality, let us consider, as a tutorial example, the one-dimensional Hubbard model subject to a finite electric field. The corresponding Hamiltonian reads

$$H = -t\sum_{\langle ij\rangle,\sigma}[e^{i\Psi(t)}c^{\dagger}_{i\sigma}c_{j\sigma} + e^{-i\Psi(t)}c^{\dagger}_{j\sigma}c_{i\sigma}] + U\sum_i n_{i\uparrow}n_{i\downarrow}, \quad (2)$$

where t is the hopping amplitude and $U > 0$ is the on-site repulsive Coulomb interaction strength. The model is exactly solvable by Bethe ansatz; see [17] and SM [33]. The real part of $\Psi(t) = Ft + i\chi$ imparts a monotonous time dependence to the excitation momenta, which on account of only being defined modulo a reciprocal lattice vector, results in the well-known Bloch oscillations as in the nondissipative case [17]. However, as we have shown above, the real part of Ψ is not responsible for renormalization of the Landau-Zener factor, which depends solely on the imaginary component. Thus in the following analysis we focus on the effect of the imaginary component of Ψ and set $\mathrm{Re}\Psi(t) = 0$. The Hubbard models with the imaginary vector potentials have been studied in the past in the context of field-driven Mott transitions [21,27]. Physically, the imaginary vector potential arises from asymmetric hopping matrix elements, respectively along and opposed to the potential gradient, in the presence of dissipation.

We now argue that the finite imaginary component of the vector potential Ψ arises naturally in any driven dissipative system endowed with the $\mathcal{PT}$-symmetric Hamiltonian. Let us express the non-Hermitian Hubbard model in Eq. (2) as a Legendre transform of a Hermitian model [28,29]:

$$H = H_0 - i\lambda J, \quad (3)$$

where H_0 is defined in Eq. (2) with $\chi = 0$, J is the current operator that commutes with H_0, and $\lambda = \sinh(\chi)$. For sufficiently small λ, the spectral gap in H_0 implies that the expectation value of J vanishes. In the opposite limit of large λ, the eigenfunctions of H are essentially those of J, and the system becomes a gapless phase with the finite steady current I. The phase transition between the equilibrium and the finite current-carrying states takes place at some critical value $|\lambda| = \lambda_c$, and the corresponding value of the Lagrange multiplier is related to the current as $I = \langle J(\lambda)\rangle$. Since H_0 and J are both Hermitian operators, for purely real λ the eigenvalues of H are real only when $\langle J(\lambda)\rangle = 0$, and complex when $\langle J(\lambda)\rangle \neq 0$. Real λ in Eq. (3) guarantees inherent $\mathcal{PT}$ symmetry of H. From the viewpoint of the $\mathcal{PT}$-symmetric quantum mechanics, the parametric region $|\lambda| < \lambda_c$ corresponds to the regime where eigenstates preserve $\mathcal{PT}$ symmetry resulting in a real spectrum for H and the zero steady current. For $|\lambda| > \lambda_c$ the spectrum of H acquires an imaginary component, and the energy gap closes leading to the finite steady current I. Thus λ_c marks the transition into the phase with broken $\mathcal{PT}$ symmetry. The $\mathcal{PT}$ symmetry breaking differs from the conventional spontaneous symmetry breaking phenomena in the sense that it is not associated with a bifurcation to degenerate ground states; rather, it manifests in the form of a complex energy spectrum and the eigenfunctions of H no longer remain eigenstates of $\mathcal{PT}$.

To see that a real λ is consistent with a nonequilibrium steady state of a dissipative driven system, we consider the dynamic equation for the density matrix ρ:

$$\frac{d\rho}{dt} = -i[H_0,\rho] - \lambda(\{J,\rho\} - 2\mathrm{tr}(\rho J)\rho). \quad (4)$$

For pure states, Eq. (4) reduces to the Schrödinger equation [30] with the non-Hermitian Hamiltonian from Eq. (3). The formal solution to Eq. (4) reads [30]

$$\rho(t) = \frac{e^{-i(H_0 - i\lambda J)t}\rho(0)e^{i(H_0+i\lambda J)t}}{\mathrm{tr}(e^{-i(H_0-i\lambda J)t}\rho(0)e^{i(H_0+i\lambda J)t})}, \quad (5)$$

where $\rho(0)$ is the density matrix corresponding to the initial state. Formula (5) generalizes the unitary evolution (which corresponds to $\lambda = 0$) onto the dissipative case ($\lambda \neq 0$) and preserves the norm, $\mathrm{tr}\rho(t) = 1$ so that $d\,\mathrm{tr}(\rho(t))/dt = 0$. The expectation value of any physical observable A is $\langle A(t)\rangle = \mathrm{tr}(\rho(t)A)$, so the expectation value of the current operator is

$$\frac{d\langle J\rangle}{dt} = -2\lambda(\langle J^2\rangle - \langle J\rangle^2). \tag{6}$$

Therefore, if λ is real, the system relaxes to a state with the constant current since $\langle J^2\rangle - \langle J\rangle^2 \geqslant 0$. This, in turn, means that the transition between the equilibrium zero-current and steady-current states caused by varying the Lagrange multiplier λ mirrors the electric field-driven dynamic Mott transition: below some critical field F_c (corresponding to λ_c) the spectral gap is finite and $I = 0$, while at larger applied fields a finite current flows with a magnitude that increases monotonically with F (or λ).

Now we complete the derivation of the electric field-induced suppression of the Landau-Zener parameter $\gamma(\chi)$ and its ultimate vanishing at $\chi = \chi_c$. Provided that $\Delta(\chi)$ is a continuous function of χ, the energy gap vanishes smoothly as $\chi \to \chi_c$. The fermionic and bosonic Mott insulator systems we study all exhibit a power-law collapse for the spectral gap, $\Delta(\chi) \sim (\chi_c - \chi)^\alpha$, with $\alpha = 0.5$ and $\alpha = 0.78 \pm 0.03$ respectively for the 1D and 2D half-filled Hubbard models, and $\alpha = 0.5$ for the (bosonic) vortex Mott insulator. This presents one of the main results of this Rapid Communication. Similar field-induced gap renormalizations have been reported in studies of other non-Hermitian models, with $\alpha = 0.5$ for the non-Hermitian XXZ chain [31] experiencing an Ising transition. Dynamic phase transition of the same universality class has also been predicted for a quasi-1D superconducting wire subject to imaginary vector potential destroying superconducting fluctuations beyond a certain threshold [2,9].

Fermionic dynamic Mott transition. To derive critical behavior of the dynamic Mott transition, we consider the Hamiltonian of the half-filled fermionic Hubbard chain [Eq. (2)], retaining only the imaginary component χ of the vector potential. Starting from the Bethe ansatz solution [21] of Eq. (2), one obtains the following expression for the Mott gap [32] on the insulating side of the transition:

$$\Delta(b) = 4t\left[u - \cosh(b) + \int_{-\infty}^{\infty}\frac{d\omega}{2\pi}\frac{J_1(\omega)e^{\omega\,\sinh(b)}}{\omega(1+2^{2u|\omega|})}\right], \tag{7}$$

where $u \equiv U/4t$, J_1 is the Bessel function of order 1, and $b \leqslant b_c$ is the parameter controlling the integration contour for the spin distribution function in the complex plane (see Ref. [21]), such that $b_c \equiv \mathrm{arcsinh}(u)$ defines the point of Mott transition, $\Delta(b_c) = 0$. We find that as $\chi \to \chi_c$, $\chi_c - \chi \approx C_1(b - b_c)^2$, which in turn leads to $\Delta(\chi) \approx C_2\sqrt{\chi_c - \chi}$, where $C_{1,2}$ are constants. Assuming that $F(\chi)$ is a well-behaved function at the threshold, we have $\Delta(F) \sim \sqrt{F_c - F}$, and we get

$$\gamma \sim (F_c - F)^{3/2} \tag{8}$$

for the scaling of the LZS tunneling parameter.

Now we consider a 2D Hubbard model, a half-filled square lattice with the nearest-neighbor hopping and nonequilibrium drive along the x direction:

$$H = \sum_{\mathbf{k},\sigma}\{-t[2(\cos(k_x) + \cos(k_y)) - i\lambda\,\sin(k_x)] - \mu\}c^\dagger_{\mathbf{k}\sigma}c_{\mathbf{k}\sigma} + U\sum_i n_{i\uparrow}n_{i\downarrow}, \tag{9}$$

where $\lambda \in [0,2]$ is the real-valued Lagrange multiplier introducing the current constraint. In what follows we calculate the spectral functions by employing a second-order perturbation theory approximation, namely the IPT within the DMFT framework [25] [see Supplemental Material (SM) [33] for details]. For a sufficiently small driving field λ and large interaction strength U, a Mott gap is formed at the Fermi level ($\omega = 0$), as seen from the spectral function in the lower inset of Fig. 2. Gradual increase of λ diminishes the gap and eventually, for $\lambda \geqslant \lambda_c$, a quasiparticle peak appears at the Fermi level by closing the gap and, hence, signifying an insulator-to-metal transition. In the upper inset of Fig. 2, we present the Mott gap Δ (extracted from the calculated spectral function) for different values of U as a function of λ. Plotting the same data as function of $(\lambda - \lambda_c(U))^\alpha$ we find them to collapse to a straight line for $\alpha = 0.78$; see the main panel. From such collapse of the data we infer a universal (i.e., independent of U) power-law behavior $\Delta \sim (\lambda - \lambda_c)^\alpha$ with $\alpha = 0.78 \pm 0.03$. This value is larger than $\alpha = 0.5$ that we obtained above for the 1D case and lies closer to $\alpha = 1$ reported in mean-field studies

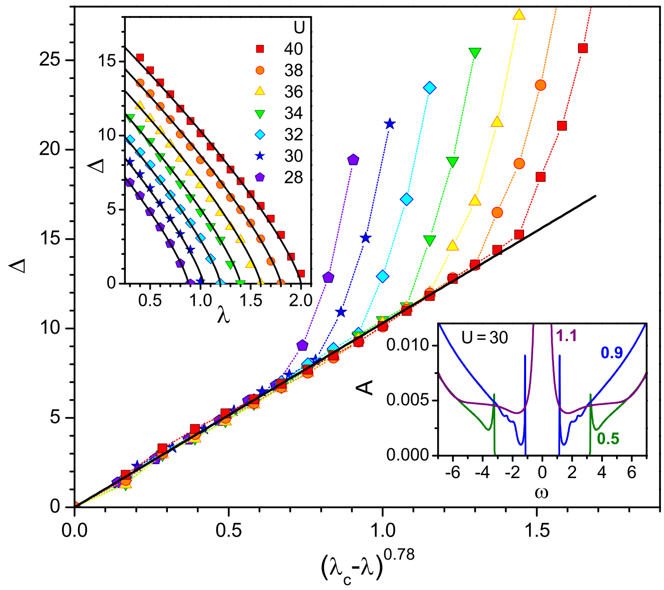

FIG. 2. Universal scaling of the Mott gap, Δ, as a function of the drive, λ, near the $\mathcal{PT}$ symmetry breaking points λ_c for different values of the Coulomb repulsion U. The upper inset shows the evolutions of the Mott gaps with increasing λ for various U, and the solid lines are fits to power laws of the form $\Delta = C(\lambda_c - \lambda)^{0.78}$. The main figure shows the same data plotted as function of $(\lambda_c - \lambda)^{0.78}$ (the legends for the inset and the main panel are the same). One sees a remarkable linear collapse of the data persisting over large region of λ. The lower inset shows the λ dependence of the spectral function $A(\omega)$ for single-particle excitations for $U = 30$. As λ increases, the spectral gap gradually narrows, and for $\lambda = 1.1$ a quasiparticle band is evident signifying a (bad) metallic phase.

of the equilibrium 2D Mott transition [34,35]. Approaching the transition from the metallic side, we also find that the effective electron mass diverges at λ_c, as is the case in the equilibrium Mott transition (see SM [33]).

Bosonic dynamic Mott transition. Let us now consider a bosonic system described by the nonrelativistic Landau-Ginzburg-Wilson (LGW) field theory. As an example, we take vortices in an array of traps in magnetic fields corresponding to integer fillings f_c. Vortex Mott insulator state has been predicted in [36] and observed in numerous experiments; see Refs. [37,38] and references therein. Moreover, we can now justly believe that the vortex Mott insulator has been seen, albeit not recognized, in numerous studies of vortex matching effects in nanopatterned superconductors [39,40]. The critical scaling behavior as a function of the driving current, temperature, and departure from commensuration $|f - f_c|$ at the dynamic vortex Mott transition has been observed in square proximity superconducting arrays [16].

Near $f_c = 1$, the vortex Mott transition is mean-field like, and its dynamics can be described by a nonrelativistic LGW effective action in Euclidean time,

$$S = \int d^2x\, d\tau \left[\psi^\dagger \frac{\partial}{\partial \tau} \psi + D|\nabla\psi|^2 + m^2|\psi|^2 + u|\psi|^4 \right]. \tag{10}$$

Here ψ is the vortex field, D is the vortex stiffness, and m and u are, respectively, the mass and the interaction parameter that govern the mean-field transition, where the "superfluid" phase of vortices corresponds to $m^2 < 0$. The applied electric current, I, exerts Magnus force on magnetic vortices. We incorporate the current into the vector potential: $A_x = It$, $A_y = 0$. We consider vortex motion in a dissipative region surrounded by a superconducting shell. This enables us to impose the simple boundary condition $\psi = 0$ outside the dissipative region. The motion in the dissipative environment is overdamped; thus the time evolution is governed entirely by Brownian processes and we can neglect Berry phase effects [first term in Eq. (10)]. In real time, the equation of motion is

$$\frac{\partial \psi}{\partial t} + \nu \frac{\delta H}{\delta \psi^*} = 0, \tag{11}$$

where $H = \int d^2x\, [D|\nabla\psi|^2 + m^2|\psi|^2 + u|\psi|^4]$ is the Hamiltonian corresponding to Eq. (10), and ν represents viscous damping of the vortex motion and is proportional to the normal resistance of the superconductor just above superconducting transition temperature T_c. Following Refs. [2,9], we ignore the nonlinear term in the vicinity of the transition and after a straightforward calculation arrive at the following expression for the distance between the two lowest-energy levels on the vortex Mott-insulating side ($I \leqslant I_c$):

$$E_1 - E_0 \approx 2E_T \sqrt{\eta\left(1 - I^2/I_c^2\right)}, \tag{12}$$

where for a size L of the normal region where the vortex is confined, $E_T = D/L^2$ is the Thouless energy, $\eta \approx (\pi^2/\sqrt{2})(I_c L/E_T \rho)$.

Interpreting the vortex excitation gap as the Mott gap, we obtain the square-root scaling behavior near the transition, and from the Landau-Dykhne formula find the barrier for vortex thermal activation to scale as

$$\gamma \sim (I_c - I)^{3/2}. \tag{13}$$

This result is in an excellent agreement with the experimental findings of Ref. [16] demonstrating scaling of the dynamic resistance near the vortex Mott transition as a function of $|I_c - I|^{3/2}/|f - f_c|$.

Discussion. To summarize, we investigated Mott transitions in fermionic half-filled Hubbard models in one and two dimensions and in bosonic vortex lattice system near integer fillings. We showed that nonequilibrium steady states of such systems are described as eigenstates of non-Hermitian Hamiltonians endowed with $\mathcal{PT}$ symmetry. The field-driven Mott transition is identified as a $\mathcal{PT}$ symmetry-breaking phenomenon. We related the driving electric field and the dissipation parameter to the non-Hermitian gauge fields governing the $\mathcal{PT}$ symmetry-breaking phase transition. While the mechanism of Mott transitions in these dissipative systems, Landau-Zener-Schwinger tunneling, is also shared with nondissipative and even noninteracting quantum systems, the key qualitative difference in the dissipative case lies in the field-induced renormalization of the excitation gap. For the 1D Hubbard chain and the bosonic 2D system, we find that the spectral gap Δ and the LZS tunneling factor γ respectively scale as $\Delta \sim (F_c - F)^{1/2}$ and $\gamma \sim (F_c - F)^{3/2}$ as a function of the driving field F. This behavior is in an accord with the current vs magnetic field scaling recently observed in the vortex Mott transition in nanopatterned superconductors [16]. For the 2D Hubbard model, we perform a DMFT analysis based on an IPT approximation scheme and obtain scaling $\Delta \sim (F_c - F)^{0.78}$. The exact solvability of our 1D even time-dependent models opens a route for quantitative investigations dynamic phenomena, in particular, the Bloch oscillations in the non-Hermitian systems. An important open problem to address is the microscopic derivation of the effective non-Hermitian Hubbard models starting from the driven Hermitian system coupled to a bath. A limitation of our Bethe ansatz approach is that it is restricted to local, nonretarded Coulomb interactions. It remains to be seen how our conclusions would change if such a general Coulomb interaction is considered. Note, finally, that vortex insulator-metal transitions have also been suggested in quantum Hall systems [41], which could provide a new platform for exploring dynamic Mott transitions. Furthermore, we hypothesize that, in general, in out-of-equilibrium open quantum systems that exhibit a dynamic phase transition (from stationary equilibriumlike to strictly nonconservative dynamics), the concurrent effects of the driving field and inherent dissipation generate imaginary parts of the systems' Hamiltonians in such a way that the resulting non-Hermitian Hamiltonians retain $\mathcal{PT}$ symmetry at drives smaller than some threshold value.

Another direction to take is the study into the role of disorder in the $\mathcal{PT}$ symmetry-breaking transitions. Past research of effects of disorder, such as vortex delocalization driven by an imaginary vector potential [22] and level statistics of zero-dimensional fermionic systems [42], are restricted to noninteracting systems. The incoherent hopping transport in disordered insulators in moderately strong electric fields, which is known to be of the directed-percolation type driven by an imaginary field [43,44], will be the subject of forthcoming publication. Our work thus paves the way for the application of $\mathcal{PT}$ symmetry concepts as a general mechanism of the dynamic phase transitions in strongly correlated quantum systems.

Acknowledgments. We thank T. I. Baturina for critical reading of the manuscript and many valuable suggestions, and D. Dhar and T. V. Ramakrishnan for illuminating discussions. The work is supported by the U.S. Department of Energy, Office of Science, Materials Sciences and Engineering Division (V.T. is partially supported by Materials Theory Institute at ANL), by the University of Chicago Center in Delhi, and DST (India) Swarnajayanti grant (No. DST/SJF/PSA-0212012-13). H.B. is grateful for support from DAE (India) and computational resources from the Department of Theoretical Physics, Tata Institute of Fundamental Research.

[1] C. M. Bender and S. Boettcher, Phys. Rev. Lett. **80**, 5243 (1998).
[2] J. Rubinstein, P. Sternberg, and Q. Ma, Phys. Rev. Lett. **99**, 167003 (2007).
[3] E. M. Graefe, H. J. Korsch, and A. E. Niederle, Phys. Rev. Lett. **101**, 150408 (2008).
[4] T. Prosen, Phys. Rev. Lett. **109**, 090404 (2012).
[5] M. J. Bhaseen, B. Doyon, A. Lucas, and K. Schalm, Nat. Phys. **11**, 509 (2015).
[6] N. Bender, S. Factor, J. D. Bodyfelt, H. Ramezani, D. N. Christodoulides, F. M. Ellis, and T. Kottos, Phys. Rev. Lett. **110**, 234101 (2013).
[7] J. Schindler, A. Li, M. C. Zheng, F. M. Ellis, and T. Kottos, Phys. Rev. A **84**, 040101(R) (2011).
[8] B. Peng, Ş. K. Özdemir, F. Lei, F. Monifi, M. Gianfreda, G. L. Long, S. Fan, F. Nori, C. M. Bender, and L. Yang, Nat. Phys. **10**, 394 (2014).
[9] N. M. Chtchelkatchev, A. A. Golubov, T. I. Baturina, and V. M. Vinokur, Phys. Rev. Lett. **109**, 150405 (2012).
[10] C. H. Ahn, J.-M. Triscone, and J. Mannhart, Nature (London) **424**, 1015 (2003).
[11] H.-T. Kim, B.-G. Chae, D.-H. Youn, G. Kim, K.-Y. Kang, S.-J. Lee, K. Kim, and Y.-S. Lim, Appl. Phys. Lett. **86**, 242101 (2005).
[12] M. S. Laad, L. Craco, and E. Müller-Hartmann, Phys. Rev. B **73**, 195120 (2006).
[13] S. Yamanouchi, Y. Taguchi, and Y. Tokura, Phys. Rev. Lett. **83**, 5555 (1999).
[14] Y. Taguchi, T. Matsumoto, and Y. Tokura, Phys. Rev. B **62**, 7015 (2000).
[15] R. Kumai, Y. Okimoto, and Y. Tokura, Science **284**, 1645 (1999).
[16] N. Poccia, T. I. Baturina, F. Coneri, C. G. Molenaar, X. R. Wang, G. Bianconi, A. Brinkman, H. Hilgenkamp, A. A. Golubov, and V. M. Vinokur, Science **349**, 1202 (2015).
[17] T. Oka and H. Aoki, Phys. Rev. B **81**, 033103 (2010).
[18] M. Eckstein, T. Oka, and P. Werner, Phys. Rev. Lett. **105**, 146404 (2010).
[19] A. G. Green and S. L. Sondhi, Phys. Rev. Lett. **95**, 267001 (2005).
[20] C. Aron, Phys. Rev. B **86**, 085127 (2012).
[21] T. Fukui and N. Kawakami, Phys. Rev. B **58**, 16051 (1998).
[22] N. Hatano and D. R. Nelson, Phys. Rev. Lett. **77**, 570 (1996); Phys. Rev. B **56**, 8651 (1997).
[23] F. Gelis and N. Tanji, Phys. Rev. D **87**, 125035 (2013).
[24] E. Shimshoni and A. Stern, Phys. Rev. B **47**, 9523 (1993).
[25] H. Barman and N. S. Vidhyadhiraja, Int. J. Mod. Phys. B **25**, 2461 (2011).
[26] A. M. Dykhne, JETP **14**, 941 (1962).
[27] R. A. Lehrer and D. R. Nelson, Phys. Rev. B **58**, 12385 (1998).
[28] T. Antal, Z. Rácz, and L. Sasvári, Phys. Rev. Lett. **78**, 167 (1997).
[29] J. Cardy and P. Suranyi, Nucl. Phys. B **565**, 487 (2000).
[30] D. C. Brody and E. M. Graefe, Phys. Rev. Lett. **109**, 230405 (2012).
[31] G. Albertini, S. R. Dahmen, and B. Wehefritz, J. Phys. A: Math. Gen. **29**, L369 (1996).
[32] E. H. Lieb and F. Y. Wu, Phys. Rev. Lett. **20**, 1445 (1968).
[33] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevB.94.041104 for details of DMFT calculations and additional results.
[34] S. Florens and A. Georges, Phys. Rev. B **70**, 035114 (2004).
[35] W. F. Brinkman and T. M. Rice, Phys. Rev. B **2**, 4302 (1970).
[36] D. R. Nelson and V. M. Vinokur, Phys. Rev. B **48**, 13060 (1993).
[37] K. Harada, O. Kamimura, H. Kasai, T. Matsuda, A. Tonomura, and V. V. Moshchalkov, Science **274**, 1167 (1996).
[38] S. Goldberg, Y. Segev, Y. Myasoedov, I. Gutman, N. Avraham, M. Rappaport, E. Zeldov, T. Tamegai, C. W. Hicks, and K. A. Moler, Phys. Rev. B **79**, 064523 (2009).
[39] A. D. Thakur, S. Ooi, S. P. Chockalingam, J. Jesudasan, P. Raychaudhuri, and K. Hirata, Appl. Phys. Lett. **94**, 262501 (2009).
[40] T. I. Baturina, V. M. Vinokur, A. Yu. Mironov, N. M. Chtchelkatchev, D. A. Nasimov, and A. V. Latyshev, Europhys. Lett. **93**, 47002 (2010).
[41] D. Shahar, D. C. Tsui, M. Shayegan, E. Shimshoni, and S. L. Sondhi, Science **274**, 589 (1996).
[42] K. B. Efetov, Phys. Rev. Lett. **79**, 491 (1997).
[43] S. P. Obukhov, Physica A **101**, 145 (1980).
[44] J. L. Cardy and R. L. Sugar, J. Phys. A: Math. Gen. **13**, L423 (1980).