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Parity-time symmetry breaking in magnetic systems
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I. INTRODUCTION

A seminal idea of parity-time ($\mathcal{PT}$)-symmetric quantum mechanics [1, 2], that has stated that the condition of Hermiticity in standard quantum mechanics required for physical observables and energy spectrum to be real can be replaced by less restrictive requirement of invariance under combined parity and time-reversal symmetry, triggered an explosive development of a new branch of science. The interpretation of $\mathcal{PT}$ symmetry as “balanced loss and gain” [3] connected $\mathcal{PT}$ symmetry-breaking to transitions between stationary and non-stationary dynamics and established its importance to understanding of the applied field-driven instabilities. Experiments on diverse variety of strongly correlated systems and phenomena including optics and photonics [4–10], superconductivity [11–13], Bose-Einstein condensates [14], nuclear magnetic resonance quantum systems [15], and coupled electronic and mechanical oscillators [16–18] revealed $\mathcal{PT}$ symmetry-breaking transitions driven by applied fields. These observations stimulated theoretical focus on far-from-equilibrium instabilities of many-body systems [12, 13, and 19] that are yet not thoroughly understood.

Here we demonstrate that non-Hermitian extension of classical Hamiltonian formalism provides quantitative description of dissipative dynamics and dynamic phase transitions in out-of-equilibrium systems. Focusing on the case of spin systems, we consider the zero-temperature spin dynamics under the action of basic non-conservative forces: phenomenological Gilbert damping [20] and Slonczewski spin-transfer torque [20] (STT). The latter serves as the most versatile way of directly manipulating magnetic textures by external currents. We propose a general complex spin Hamiltonian, in which non-Hermiticity reflects dissipation and deviation from equilibrium. The imaginary part of the proposed spin Hamiltonian describes the expectation value of the Hamiltonian proves extremely useful for the general understanding of the applied field-driven instabilities. Experiments on diverse variety of strongly correlated systems and phenomena ranging from superconductivity to cold-atom and two-level systems, our results provide quantitative perspectives on the nature of phase transitions associated with $\mathcal{PT}$ symmetry-breaking in a broad class of far-from-equilibrium systems.

We introduce the non-Hermitian Hamiltonian for a single spin operator $\hat{S}$:

$$\hat{H} = \frac{E(\hat{S}) + i \hat{J} \cdot \hat{S}}{1 - i \alpha},$$

where $E(\hat{S})$ denotes the standard Hermitian spin Hamiltonian determined by the applied magnetic field $\vec{H}$ and magnetic anisotropy constants $k_i$ in the $x,y,z$-directions: $E(\hat{S}) = \sum_i k_i \vec{S}_i^2 + \gamma \vec{H} \cdot \vec{S}$. A schematic system setup is shown in Fig. 1. The phenomenological constant $\alpha > 0$ in Eq. (1) describes damping; the imaginary field $i\alpha$ is responsible for the applied Slonczewski STT, with $\hat{J} = \hat{e}_p (\hbar/2e)\eta J$ being the spin-angular momentum deposited per second in the direction $\vec{e}_p$ with spin-polarization $\eta = (J_1 - J_2)/(J_1 + J_2)$ of the incident current $J$; $\gamma = g\mu_B / h$ is the absolute value of the gyromagnetic ratio; $g \approx 2$, $\mu_B$ is the Bohr magneton; $h$ is the Planck’s constant, and $e$ is the elementary charge. We conjecture that Eq. (1) serves as a fundamental generalization of the Hamiltonian description of both quantum and classical spin systems, which constitutes one of our core results. This form of the Hamiltonian proves extremely useful for the general understanding of STT-driven dissipative spin dynamics. In this work we focus primarily on the classical limit of spin dynamics, while the semiclassical limit of finite spin will be considered elsewhere.

Spin dynamics in the classical limit is conveniently obtained by studying expectation value of the Hamiltonian (1) with respect to SU(2) spin-coherent states [23, 24]: $|z\rangle = e^{i\hat{S}_z}|\hat{S}, -\hat{S}\rangle$, where $\hat{S}_x \equiv \hat{S}_+ + i\hat{S}_- $, and $z \in \mathbb{C}$ is the standard stereographic projection of the spin direction on a unit
sphere, \( z = (s_z + is_y)/(1 - s_x) \), with \( s_i \equiv S_i/S \). Note that such parametrization of the phase space for a classical single spin system (i.e. in the limit \( S \to \infty \)) guarantees the invariance of the traditional equation of motion [24] under generalization to non-Hermitian Hamiltonians (see Supplemental Material):

\[
\dot{z} = i\left(1 + \bar{z}z\right)^2 \frac{\partial\mathcal{H}}{\partial\bar{z}},
\]

(2)

where \( z \) and \( \bar{z} \) form a complex conjugate pair of stereographic projection coordinates, and

\[
\mathcal{H}(z, \bar{z}) = \left\langle \frac{z\bar{\mathcal{H}}(z)}{z\bar{z}} \right\rangle
\]

(3)

is the expectation value of the Hamiltonian (1) in spin-coherent states (for a detailed review see, e.g., Ref. [25]). In this formulation, the eigenstates of \( \hat{\mathcal{H}} \) correspond to the fixed points \( z_i \) of the equation of motion for \( \mathcal{H} \), while the eigenvalues (i.e. energy values) are equal to \( \mathcal{H} \) evaluated at the corresponding fixed points, \( E_i = \mathcal{H}(z_i, \bar{z}_i) \).

Assuming a constant magnitude of the total spin, \( \dot{S} = 0 \), Eq. (2) reduces to the following equation of spin dynamics in the classical limit:

\[
\dot{S} = \nabla_S (\text{Re}\mathcal{H}) \times S + \frac{1}{S} \left[ \nabla_S (\text{Im}\mathcal{H}) \times S \right] \times S.
\]

(4)

Here we refer to the real and imaginary parts of the Hamiltonian function \( \mathcal{H} \) written in the spin-S-representation. For the non-Hermitian Hamiltonian (1), Eq. (4) reproduces the LLGS equation describing dissipative STT-driven dynamics of a macrospin:

\[
\begin{align*}
(1 + \alpha^2)\dot{S} &= \gamma H_{\text{eff}} \times S + \frac{\alpha \gamma}{S} [\gamma H_{\text{eff}} \times S] \times S + \frac{1}{S} S \times [S \times \mathbf{j}] \\
&+ \alpha S \times \mathbf{j}, \\
\gamma H_{\text{eff}} &= \nabla_S E(S).
\end{align*}
\]

(5)

The first two terms in Eq. (5) describe the standard Landau–Lifshitz torque and dissipation, while the last two are responsible for the dissipative (‘anti-damping’) and conservative (‘effective field’) Slonczewski STT contributions, correspondingly, both of which appear naturally from the imaginary magnetic field term in the Hamiltonian (1).

**\( \mathcal{PT} \)-Symmetric Hamiltonian**

Slonczewski STT turns the total spin-angular momentum, \( S \), in the direction of spin-current polarization, \( \mathbf{e}_p \), without changing its magnitude. On the \( S \)-sphere this can be represented by a vector field converging in the direction of \( \mathbf{e}_p \) and originating from the antipodal point. It is the imaginary magnetic field \( \mathbf{j} \) that produces exactly the same effect on spin dynamics, according to Eq. (2). The action of STT is invariant under the simultaneous operations of time-reversal and reflection with respect to the direction \( \mathbf{e}_p \), which is the underlying reason behind the inherent \( \mathcal{PT} \)-symmetry of certain STT-driven magnetic systems, including the one considered below.

Before turning to the \( \mathcal{PT} \)-symmetric form of Hamiltonian (1), we note that \( \mathcal{PT} \)-symmetric systems play an important role in the studies of non-equilibrium phenomena and provide a unique non-perturbative tool for examining the phase transition between stationary and non-stationary out-of-equilibrium dynamics. We show that despite being non-Hermitian, such systems can exhibit both of the above types of behavior, depending on the magnitude of the external non-conservative force. In the parametric regime of unbroken \( \mathcal{PT} \)-symmetry, systems exhibit physical properties seemingly equivalent [26] to those of Hermitian systems: real energy spectrum, existence of integrals of motion (see Supplemental Material), and, notably, the validity of the quantum Jarzynski equality [27]. However, in the regime of broken \( \mathcal{PT} \)-symmetry, one observes complex energy spectrum and non-conservative dynamics. Therefore, the ‘true’ transition between stationary and non-stationary dynamics can be identified as the \( \mathcal{PT} \)-symmetry-breaking phase transition.

Spin systems are generally subject to various non-linear magnetic fields including ones originating from shape, exchange, magnetocrystalline and magnetoelastic anisotropies. Restricting ourselves for simplicity to a second-order anisotropy term, we arrive at the following Hamiltonian for non-linear magnetic system with uniaxial anisotropy and applied Slonczewski STT:

\[
\mathcal{H}_{\mathcal{PT}} = \gamma H_0 \left( k_z \hat{S}_z^2 + h_x \hat{S}_x + i\beta \hat{S}_y \right),
\]

(7)

where the applied magnetic field \( h_x \) is measured in units of some characteristic magnetic field \( H_0 \), and \( \beta \) is a dimensionless STT parameter determining the relative to \( \dot{S} \) amount of
FIG. 2. Real (a) and imaginary (b) parts of energy spectrum of the Hamiltonian (7) as a function of the STT parameter $\beta$ for $h_z = 1$ and $D = 20$. Blue and red lines correspond to the eigenvalues $E_{1,2}$ and $E_{3,4}$, respectively. The first PT symmetry-breaking transition occurs at $|\beta| = \beta_1 \approx 4.5$.

angular momentum transferred in time $\tau \equiv (\gamma H_0)^{-1}$ (characteristic timescale of the dynamics, used as a unit of dimensionless time in what follows). The Hamiltonian (7) modeling the dynamics of the free magnetic layer in a typical nanopillar device with fixed polarizer layer (see Fig. 1) is PT-symmetric: it is invariant under simultaneous action of parity and time-reversal operators $(y \rightarrow -y, t \rightarrow -t, i \rightarrow -i)$. Because the Hamiltonian $\hat{H}_{PT}$ commutes with an anti-linear operator $\hat{P}^T$, its eigenvalues are guaranteed to appear in complex conjugate pairs. Notice that PT-symmetric Hamiltonian (7) does not contain damping, which is assumed to be negligibly small, as is the case in many experimental systems.

CLASSICAL SPIN SYSTEM

In order to best illustrate the mechanism of PT symmetry-breaking, we focus on the classical limit, $S \rightarrow \infty$ and $k, S \rightarrow D/2$, where $D$ is the dimensionless uniaxial anisotropy constant. Formula (2) then yields the following equation of motion for the Hamiltonian (7):

$$\dot{z}(t) = -\frac{i(h_z + \beta)}{2} \left( \frac{h_z - \beta}{h_z + \beta} \right) - iD z \left( 1 - \frac{1}{1 + |z|^2} \right),$$

with up to six fixed points $z_k, k = 1, \ldots, 6$.

Shown in Fig. 2 are the real and imaginary parts of the energy spectrum $E_{1-6}$ as functions of the STT amplitude $\beta$. It reveals that in a system with strong anisotropy, $D \gg 1$, PT symmetry breaking occurs in three separate transitions, with the first one at $|\beta| = \beta_1 = |h_z| \sqrt{1 + \sqrt{1 + (2D/|h_z|)^2}}/2$, which corresponds to the smallest amplitude of STT at which $\text{Im}(E) \neq 0$. Therefore, PT symmetry is not broken in the entire phase space of initial spin directions simultaneously, at variance to the linear spin system with $D = 0$ (see Supplementary Material). Instead, the regions of broken and unbroken PT symmetry may coexist in the phase diagram of a non-linear spin system.

In what follows we consider a system described by the Hamiltonian (7) with $h_z = 1$ and $D = 20$. For all $|\beta| < \beta_1 \approx 4.5$, PT symmetry is unbroken and the character of spin (magnetization) dynamics is oscillatory in the entire phase diagram, i.e. for all possible initial conditions $z$. At $|\beta| = \beta_1$ the phase transition (first of the three, see Fig. 2) occurs sharply in a wide region around the easy plane, $|z| = 1$, i.e. near the equator of the unit $S$-sphere, shown in gray in Fig. 3a, b in Cartesian and stereographic projection coordinates, correspondingly. It this region the nature of spin dynamics becomes fundamentally different—all spin trajectories follow the lines connecting the fixed points $z_1$ and $z_2$, where $z_{1,2} = -(D h_z \pm \sqrt{|\beta|^2 - |h_z|^2 - D^2 T^2})(h_z + \beta)$, and no closed trajectories are possible, see Fig. (3)b.

As $|\beta|$ is increased further, the region of broken PT symmetry expands until it eventually closes around the fixed point $z_3$ at $\beta_2 \approx 9.3$ (second bifurcation in Fig. 2) and, eventually, the last region of unbroken PT symmetry near $z_3$ disappears at $\beta_3 \approx 10.8$. The second and third phase transitions are less relevant experimentally as they occur in the vicinity of the least favorable spin directions (parallel and anti-parallel to the hard axis $z$) and at considerably higher applied currents.

The predicted transition from precessional dynamics (unbroken PT symmetry) to exponentially fast saturation in the direction $z_3(h_z, \beta)$ for any initial spin position around the easy plane (broken PT symmetry) occurs in the setup with mutually perpendicular applied magnetic field and Slonczewski STT. Such a transition in nanoscale magnetic structures can be used for STT- or magnetic field-controlled magnetization switching in spin valves and a variety of other experimental systems. This effect can further be used for direct measurements of the amplitude of the applied STT, which, unlike the applied current, can be hard to quantify experimentally.

**NUMERICAL SIMULATIONS OF PT SYMMETRY BREAKING**

Here we present the results of numerical simulations confirming the PT symmetry-breaking phase transition in the classical single spin system (7) by modeling magnetization dynamics of a ferromagnetic disk 100nm in diameter and
The permalloy disk was simulated in external magnetic field applied along the $x$-axis, $H_0 = 400$ Oe, which corresponds to the characteristic time $\tau \approx 0.14\text{ns}$. The STT was produced by applying electric current perpendicular to the disk in the $z$-direction with spin polarization $\eta = 0.7$ along $e_p = \hat{y}$ (see Fig. 1) and current density $\beta$ measured in dimensionless units of $2eH_0M_{sat}d/\eta \hbar \approx 0.7 \times 10^3 \text{A/cm}^2$. While such current density is comparable to typical switching current densities in STT-RAM devices [29, 30], its magnitude can be optimized for various practical applications by changing $H_0$ and adjusting the size, shape and material of the ferromagnetic element.

For all possible initial spin directions $z$, we calculated the critical amplitude of the applied STT, $\beta^*$, for which the character of spin dynamics changes from oscillatory (at $|\beta| < \beta^*$) to exponential saturation. Shown in Fig. 3c is the color map of $\beta^*$ as a function of $z$ in complex stereographic coordinates. The region shown in the shades of blue corresponds to the initial conditions $z$, for which the minimum values of $\beta$ that would guarantee saturation of spin dynamics in the direction of $z$, in under $0.5\text{ns}$ are between 4.6 and 4.8. This is in full agreement with the region of broken $\mathcal{PT}$ symmetry at $\beta = 4.7$ calculated analytically, i.e. the shaded area in Fig. 3b (the outline is repeated in Fig. 3c for comparison). Outside of this region, a considerably larger magnitude of the applied STT is required to break $\mathcal{PT}$ symmetry.

The agreement between theoretical results and micromagnetic simulations is remarkable considering the non-zero Gilbert damping parameter ($\alpha = 0.01$) and non-linear effects (demagnetizing field, finite size/boundary effects, etc.) inherently present in the micromagnetic simulations but not included in the model Hamiltonian (7).

**CONCLUSION**

The presented non-Hermitian Hamiltonian formulation of dissipative non-equilibrium spin dynamics generalizes the previous result [31], where the classical Landau–Lifshitz equation was derived from a non-Hermitian Hamilton operator, to open STT-driven spin systems. The introduction of Slonczewski STT in the imaginary part of the Hamiltonian revealed the possibility of STT-driven $\mathcal{PT}$ symmetry-breaking phase transition. Micromagnetic simulations confirm the $\mathcal{PT}$ symmetry-breaking phenomenon in realistic mesoscopic magnetic systems and its robustness against weak dissipation, indicating high potential for impacting spin-based information technology. The way STT enters the complex Hamiltonian (1), i.e. as imaginary magnetic field, provides a unique tool for studying Lee-Yang zeros [32] in ferromagnetic Ising and Heisenberg models and, more generally, dynamics and thermodynamics in the complex plane of physical parameters. We envision further realizations of the $\mathcal{PT}$ symmetry-breaking phase transitions in diverse many-body condensed-matter systems and the expansion of practical implementations of the $\mathcal{PT}$ symmetry beyond the present realm of optics [33] and acoustics [34].

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