

## Quantum spin transfer torque induced nonclassical magnetization dynamics and electron-magnetization entanglement

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The standard spin transfer torque (STT)—where spin-polarized current drives the dynamics of magnetization viewed as a classical vector—requires noncollinearity between electron spins carried by the current and magnetization of a ferromagnetic layer. However, recent experiments [A. Zholud *et al.*, *Phys. Rev. Lett.* **119**, 257201 (2017)] observing magnetization dynamics in spin valves at cryogenic temperatures, even when electron spin is collinear to magnetization, point at overlooked quantum effects in STT that can lead to *highly nonclassical* magnetization states. Using quantum many-body treatment, where an electron injected as a spin-polarized wave packet interacts with local spins comprising the anisotropic quantum Heisenberg ferromagnetic chain, we define quantum STT as any time evolution of local spins due to an initial many-body quantum state not being an eigenstate of an electron+local-spins composite system. For time evolution caused by injected spin- $\downarrow$  electron scattering off local  $\uparrow$ -spins, *entanglement* between electron and local spin subsystems takes place leading to *decoherence* and, therefore, shrinking of the total magnetization but without rotation from its initial orientation, which is compatible with the experiments. Furthermore, the same processes—entanglement and thereby induced true decoherence—are present even in the standard noncollinear geometry, intertwined with the usual magnetization rotation. This is because STT in a quantum many-body picture is always caused by an electron spin- $\downarrow$  factor state, and the only difference between collinear and noncollinear geometries is in the relative size of the contribution of the initial separable state containing such a factor to the superposition of separable many-body quantum states generated by time evolution.

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The standard spin transfer torque (STT) [1], predicted in the seminal works of Slonczewski [2] and Berger [3], is a phenomenon where a flux of spin-polarized electrons injected into a ferromagnetic metal (FM) layer drives its magnetization dynamics. The origin of STT is transfer of spin angular momentum from electrons to local magnetic moments of the FM layer, so it is fundamentally a nonequilibrium quantum many-body physics effect. Nevertheless, local magnetic moments  $\mu_S \mathbf{S}(\mathbf{r})$  are typically treated as *classical vectors of fixed length* [1,4] whose dynamics is governed by the Landau-Lifshitz-Gilbert (LLG) equation [5] extended by adding the STT term [6–8]

$$\mathbf{T} \propto \langle \hat{\mathbf{s}}_e \rangle \times \mathbf{S}(\mathbf{r}). \quad (1)$$

Thus, the nonequilibrium spin density  $\langle \hat{\mathbf{s}}_e \rangle$  caused by the flowing electrons *must be noncollinear* to the direction of local spin  $\mathbf{S}(\mathbf{r})$  to drive magnetization dynamics in such a classical picture. The dynamics can include oscillations or complete reversal, whose conversion into resistance variations has emerged as a key resource for next-generation spintronic technologies, such as nonvolatile magnetic random access memories, microwave oscillators, microwave detectors, spin-wave emitters, memristors, and artificial neural networks [9–11].

For example, passing current through a spin valve tri-layer fixed-FM/normal-metal/free-FM, as employed in early experiments on standard STT [12,13], first causes the FM layer with fixed magnetization to spin-polarize the current, which then impinges onto the second FM layer with free magnetization that fluctuates in the classical picture due to a random magnetic field caused by thermal motion. When impinging spins and fluctuating magnetization become noncollinear, standard STT can either amplify such fluctuations (for the fixed-to-free spin current direction) or reduce them (for the free-to-fixed spin current direction), as predicted theoretically [14] and confirmed experimentally [15] at room temperature.

However, this well-established picture *cannot* explain very recent experiments [16] on collinear spin valves at cryogenic temperatures  $\lesssim 3$  K, where resistance measurements have revealed magnetization dynamics even though thermal fluctuations that could introduce noncollinearity between the free and fixed magnetizations are suppressed. This implies a mechanism where standard STT is zero,  $\mathbf{T} \equiv \mathbf{0}$  in Eq. (1), so that magnetization does not rotate from the initial configuration. Nevertheless, it changes its length, thereby signaling generation of *highly nonclassical* magnetization states [16]. However, the proposed intuitive picture [16] where such a mechanism would amplify quantum spin fluctuations, for both fixed-to-free and free-to-fixed spin current directions, cannot be rigorously justified. That is, although quantum fluctuations of the local spin operators [17] (or, equivalently, zero-point

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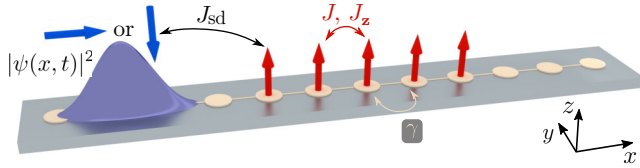


FIG. 1. Schematic view of a quantum many-body system consisting of a FM layer whose  $N = 5$  local spins comprise the open  $XXZ$  quantum Heisenberg ferromagnetic chain with anisotropic exchange interactions  $J_x > J$ , and they are attached to a 1D TB chain (composed of  $L_x = 400$  sites) where electron hops with parameter  $\gamma$ . The spin-polarized electron wave packet is injected along the TB chain, with its spin pointing in the  $-z$  or  $+x$  direction, which is collinear [a case in which the standard STT of Slonczewski [2] and Berger [3] in Eq. (1) is absent] or noncollinear, respectively, to local spins pointing initially along the  $+z$  direction. The spin of the wave packet interacts with local spins via  $s$ - $d$  exchange interaction  $J_{sd}$ .

energy of magnons as bosonic particles to which spin operators can be mapped) play an important role in lowering the energy of classical ground states of antiferromagnets [18] or noncollinear spin textures [19], they *vanish* in a FM with uniaxial anisotropy because the collinear state of local magnetic moments is also a ground eigenstate of the exact Hamiltonian [20].

Aside from a few disparate attempts [21–23], a general framework for describing current-driven *quantum dynamics* of magnetization is lacking. Note that quantum transport theories, such as the nonequilibrium Green function formalism [7,8,24,25] or the scattering matrix approach [26,27], are routinely used to compute  $\langle \hat{s}_e \rangle$  in Eq. (1) for a given single-particle Hamiltonian, but this serves only as an input [7,8] for the LLG equation describing classical dynamics of magnetization. The LLG equation can be justified under the assumptions [5] of large spin  $S \rightarrow \infty$ ,  $\hbar \rightarrow 0$  (while  $S \times \hbar \rightarrow 1$ ), and in the *absence of entanglement*. The latter assumption means that local spins comprising the total magnetization should remain in a separable quantum state,  $|S_1\rangle \otimes |S_2\rangle \otimes \dots \otimes |S_N\rangle$ , as exemplified by the ground state of a FM,  $|\uparrow\rangle \otimes |\uparrow\rangle \otimes \dots \otimes |\uparrow\rangle$ .

Instead of classical micromagnetics [4,14] or a quantum-classical [7,8] description of standard-STT-induced magnetization dynamics, here we introduce a quantum many-body picture of *both* flowing-electron-spin–local-spins interactions and the ensuing time evolution of local spins at zero temperature. For this purpose, we employ a system depicted in Fig. 1 where a spin-polarized electron wave packet, assumed to originate from a fixed FM layer, is injected along a one-dimensional (1D) tight-binding (TB) chain whose sites in the middle host local spins comprising a quantum Heisenberg ferromagnetic chain modeling the free FM layer. The states of such a composite quantum system electron+local-spins reside in the Hilbert space

$$\mathcal{H} = \mathcal{H}_{\text{orb}}^e \otimes \mathcal{H}_{\text{spin}}^e \otimes \mathcal{H}_{\text{spin}}^1 \cdots \otimes \mathcal{H}_{\text{spin}}^N, \quad (2)$$

which is the tensor product of the orbital electron subspace  $\mathcal{H}_{\text{orb}}^e$  (of finite dimension equal to the number of sites  $L_x$  of the TB chain), the two-dimensional subspace  $\mathcal{H}_{\text{spin}}^e$  for electron spin, and  $\mathcal{H}_{\text{spin}}^n$  as two-dimensional subspaces for

$n = 1, \dots, N$  local spins assumed to be spin- $\frac{1}{2}$  as well. The system Hamiltonian acting in  $\mathcal{H}$  is

$$\hat{H} = -\gamma \sum_{\langle ij \rangle} |i\rangle \langle j| - J_{sd} \sum_i |i\rangle \langle i| \otimes \hat{s}_e \cdot \hat{S}_i(t) - \sum_{\langle ij \rangle} [J(\hat{S}_i^x \cdot \hat{S}_j^x + \hat{S}_i^y \cdot \hat{S}_j^y) + J_z \hat{S}_i^z \cdot \hat{S}_j^z], \quad (3)$$

where  $|i\rangle$  is the electron orbital centered on site  $i$ ,  $\gamma = 1$  eV is hopping between nearest-neighbor sites, and  $J_{sd} = 0.1$  eV is the strength of  $s$ - $d$  exchange interaction between electron and local spins. The exchange interaction between the nearest-neighbor local spins is  $J = 0.1$  eV and  $J_z = 0.1005$  eV, which are slightly different in order to include the uniaxial anisotropy, quantified by the parameter  $\Delta = J_z/J$  [28,29], with the  $z$  axis as the easy axis. The third term in Eq. (3) is denoted as the  $XXZ$  quantum Heisenberg ferromagnetic chain with *open* boundary conditions (due to the first and last spin interacting with only one nearest-neighbor spin) [28,29]. The spin operators in Eq. (3) are constructed as  $\hat{s}_e = \hat{I} \otimes \hat{\sigma} \otimes \hat{I} \otimes \dots \otimes \hat{I}$  for electron spin,  $\hat{S}_1 = \hat{I} \otimes \hat{I} \otimes \hat{\sigma} \otimes \hat{I} \otimes \dots \otimes \hat{I}$  for first local spins, and analogously for all other local spins, where  $\hat{\sigma} = (\hat{\sigma}_x, \hat{\sigma}_y, \hat{\sigma}_z)$  is the vector of the Pauli matrices and  $\hat{I}$  is the unit operator. The eigenspectrum of an isolated  $XXZ$  chain is shown in Fig. 2(a), while the eigenspectrum of the whole many-body Hamiltonian in Eq. (3) is shown in Fig. 2(b). The ground state in the former (latter) case has degeneracy six (seven), as shown in Fig. 2(c) and expected for a system of coupled five (six) spin- $\frac{1}{2}$  angular momenta.

At  $t = 0$ , the many-body quantum state is a separable one,

$$\langle x | \Psi(t=0) \rangle = C e^{ik_x x - \delta k_x^2 x^2 / 4} \otimes \chi_e \otimes \chi_1 \otimes \dots \otimes \chi_N. \quad (4)$$

Its first factor in  $\mathcal{H}_{\text{orb}}^e$  is chosen as a Gaussian wave packet with momentum along the  $+x$  direction and centered on the left edge of the TB chain, as illustrated in Fig. 1, where  $C$  is the normalization constant. To mimic the current of electrons at the Fermi level that interact with the ground state of the free FM layer within a spin valve, we use  $k_x a = 0.1$  and  $\delta k_x a = 0.2$  ( $a$  is the lattice spacing), which tune the wave-packet average energy  $E = -2.36$  eV and its standard deviation  $\delta E = 0.054$  eV to be close to the ground-state eigenenergy  $E_0 = -2.43$  eV [Fig. 2(b)] of the Hamiltonian in Eq. (3). In the ground state, all local spins are aligned with the anisotropy  $z$  axis, as shown in Fig. 2(c), so we choose  $\chi_n = \begin{pmatrix} 1 \\ 0 \end{pmatrix}$  for  $n = 1, \dots, N$ . To mimic *minority*-spin electrons impinging onto the free FM layer within a spin valve with parallel magnetizations, we select initial spin polarization of the wave packet in the  $-z$  direction, as described by the spinor  $\chi_e = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$ . For the standard STT setup with noncollinear magnetizations of the fixed and free FM layers, we use spin polarization in the  $+x$  direction,  $\chi_e = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 \\ 1 \end{pmatrix}$ .

For transparency of the discussion, operating with a small number of excited states of the  $XXZ$  chain that can be analyzed one by one, we employ a small number  $N = 5$  of local spins. The chosen length  $L_x = 400$  of the TB chain ensures that the wave packet does not reflect from its boundaries within the time frame considered in Figs. 3 and 4. The *numerically exact*  $|\Psi(t)\rangle$ , governed by the Schrödinger

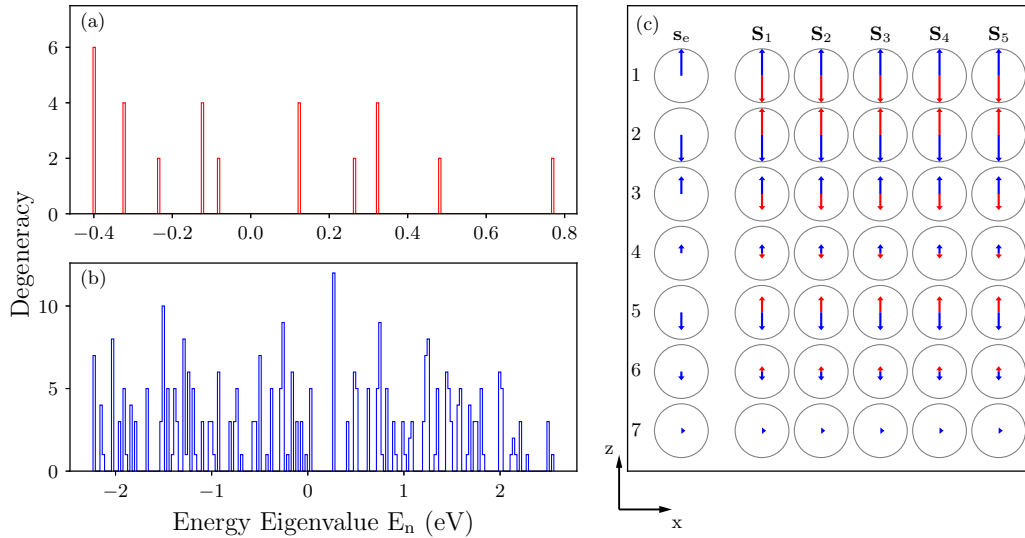


FIG. 2. (a) Eigenspectrum of the XXZ quantum Heisenberg ferromagnetic chain whose  $N = 5$  local spins in Fig. 1 do not interact with electron spin ( $J_{sd} = 0$ ). (b) Eigenspectrum of the many-body Hamiltonian in Eq. (3) whose local spins interact via  $s$ - $d$  interaction ( $J_{sd} = 0.1$  eV) with electron spin within a TB chain of length  $L_x = 400$ . (c) Expectation value of electron spin (first column) and local spins, extracted from their subsystem density matrices via Eq. (6), in the degenerate ground state of the lowest energy in panel (a) (red arrows) or (b) (blue arrows).

equation  $i\hbar\partial|\Psi(t)\rangle/\partial t = \hat{H}|\Psi(t)\rangle$ , is obtained by using the Crank-Nicolson algorithm [30].

Figure 2(c) shows that the degenerate ground state has electron and local spins parallel to each other due to  $s$ - $d$  interaction between them acting to align them. Thus, when an electron with spin- $\downarrow$  along the  $-z$  direction is injected, its spin is collinear to local spins, but  $|\Psi(t=0)\rangle \equiv |G\rangle \otimes |\downarrow_e; \uparrow_1 \cdots \uparrow_N\rangle$  (this form is used below for economy of notation) at  $t=0$  is not an eigenstate of the Hamiltonian in Eq. (3). This causes time evolution of the electron-spin+local-spins subsystem, which rigorously defines quantum STT even in a situation in which the standard STT in Eq. (1) is identically zero. Over the course of time evolution,  $|\Psi(t)\rangle$  becomes an entangled state due to linear superpositions of separable states being generated for  $t > 0$ . The entanglement entails that each subsystem must be described using the appropriate reduced density matrix [31]

$$\hat{\rho}_{\text{sub}} = \text{Tr}_{\text{other}}|\Psi(t)\rangle\langle\Psi(t)|, \quad (5)$$

obtained via partial trace applied to the pure state density matrix  $|\Psi(t)\rangle\langle\Psi(t)|$ . For example, tracing over the states in the subspace  $\mathcal{H}_{\text{orb}}^e \otimes \mathcal{H}_{\text{spin}}^e \otimes \mathcal{H}_{\text{spin}}^2 \cdots \otimes \mathcal{H}_{\text{spin}}^N$  yields the density matrix of the first local spin,

$$\hat{\rho}_1(t) = \frac{1}{2}[\hat{I} + \mathbf{S}_1(t) \cdot \hat{\boldsymbol{\sigma}}], \quad (6)$$

where  $\mathbf{S}_1(t) = \text{Tr}[\hat{\rho}_1(t)\hat{\boldsymbol{\sigma}}]$  is the spin expectation value (in units of  $\hbar/2$ ), also denoted as the polarization (or Bloch) vector [31]. Pure (or fully coherent) quantum states of spin- $\frac{1}{2}$  are characterized by  $|\mathbf{S}_1| = 1$ , while  $0 < |\mathbf{S}_1| < 1$  signifies their decoherence [31,32] toward mixed (or partially coherent [34]) states. Figure 3(c) shows that the first local spin has  $S_1^z < 1$ ,  $S_1^y = S_1^x \equiv 0$ , and, therefore,  $|\mathbf{S}_1| < 1$ . The electron spin also exhibits decoherence,  $|s_e| < 1$ , in Fig. 3(a). Virtually the same time dependences as in Fig. 3(c) are obtained for other local spins  $i = 2, \dots, 5$ , and, therefore, for total magnetization as the sum of local spins. Thus, this is precisely the

highly nonclassical state of magnetization conjectured from the measurement of the spin valve resistance [16], which increases  $\propto 1 - M_z$  due to magnetization  $M_z = g\mu_B \sum_i S_i^z$  shrinking without rotation (i.e.,  $M_x = M_y = 0$ ) away from its initial orientation.

To explain the origin of magnetization decoherence, or, equivalently, of the subsystem comprised of *all* local spins, we view the multipartite [due to  $N+2$  factors in Eq. (2)] total system as a bipartite one, i.e., as being composed of the electron subsystem whose states reside in  $\mathcal{H}_{\text{orb}}^e \otimes \mathcal{H}_{\text{spin}}^e$  and the subsystem of all local spins. The purity of the latter is defined as [31,32]

$$\mathcal{P}_{\text{spins}}^{\text{local}}(t) = \text{Tr}\{[\hat{\rho}_{\text{spins}}^{\text{local}}(t)]^2\}, \quad (7)$$

where the density matrix  $\hat{\rho}_{\text{spins}}^{\text{local}}(t)$  is obtained via Eq. (5) by tracing over the states in the subspace  $\mathcal{H}_{\text{orb}}^e \otimes \mathcal{H}_{\text{spin}}^e$ . The decay of  $\mathcal{P}_{\text{spins}}^{\text{local}}(t)$  below one in Fig. 3(b) quantifies *true decoherence* [32,33] (i.e., the decoherence that cannot be attributed to any classical fluctuations) of an initially pure state  $|\uparrow\rangle \otimes |\uparrow\rangle \otimes \cdots \otimes |\uparrow\rangle$  as the decay [31,32] of the off-diagonal elements of  $\hat{\rho}_{\text{spins}}^{\text{local}}(t)$  caused by entanglement with the electron subsystem. The purity of a decohered electron subsystem in Fig. 3(b) is identical to that of the local spin subsystem, as expected for entanglement in bipartite quantum systems [31,32].

To understand the states of the electron-spin+local-spins subsystem that are excited during time evolution initiated by injection of a single spin-polarized electron, we compute the density matrix  $\hat{\rho}_{\text{spins}}^{\text{e+local}}(t)$  of this subsystem obtained by partial trace in Eq. (5) performed over the states in  $\mathcal{H}_{\text{orb}}^e$ . The probability of finding this subsystem in state  $|\sigma_e; \sigma_1, \dots, \sigma_N\rangle$  at time  $t$ ,

$$\text{prob}_{\text{spins}}^{\text{e+local}}(t) = \langle \sigma_e; \sigma_1, \dots, \sigma_N | \hat{\rho}_{\text{spins}}^{\text{e+local}}(t) | \sigma_e; \sigma_1, \dots, \sigma_N \rangle, \quad (8)$$

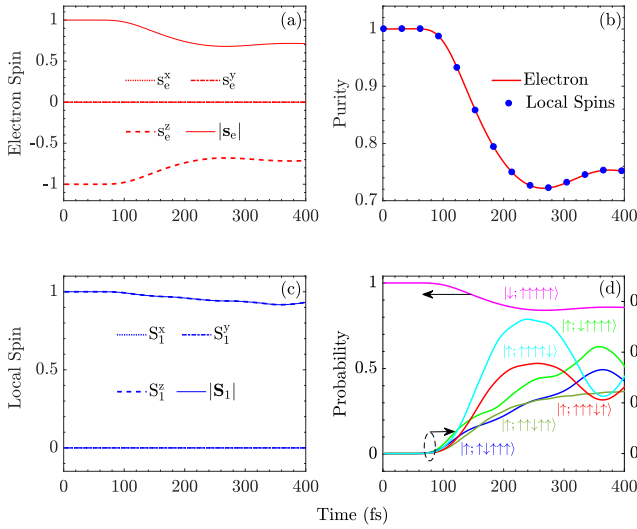


FIG. 3. Time dependence of the expectation value of spin (in units  $\hbar/2$ ) obtained from spin- $\frac{1}{2}$  density matrix in Eq. (6) for (a) spin of an injected electron wave packet in Fig. 1, which at  $t = 0$  points in the  $-z$  direction that is *collinear* and antiparallel to local spins pointing in the  $+z$  direction; and (c) first local spin in Fig. 1 [the time dependence of the expectation value of local spins  $n = 2-5$  is nearly identical to (c)]. (b) Purity defined in Eq. (7) of the subsystem composed of electron degrees of freedom (orbital and spin) or of the subsystem composed of all local spins. (d) Probability in Eq. (8) to find an electron-spin+local-spins subsystem in a many-body quantum state  $|\sigma_e; \sigma_1 \sigma_2 \sigma_3 \sigma_4 \sigma_5\rangle$ .

is shown in Fig. 3(d) for an electron injected with spin along the  $-z$  direction. The subspace of  $\mathcal{H}$  whose states can generate nonzero  $\text{prob}_{\text{spins}}^{\text{e+local}}(t)$  is restricted by the energy bands in Fig. 2(b) (caused by anisotropy and boundaries [28,29]) and symmetries, such as the fact that total spin in the  $z$  direction has to be conserved due to its operator,  $\hat{S}_{\text{tot}}^z = \hat{s}^z + \hat{S}_1^z + \dots + \hat{S}_N^z$ , commuting with the Hamiltonian in Eq. (3),  $[\hat{H}, \hat{S}_{\text{tot}}^z] = 0$ . Because of the latter requirement, all states  $|\sigma_e; \sigma_1, \dots, \sigma_N\rangle$  participating in time evolution must have the same number of  $\uparrow$ -spins, so that one finds in Fig. 3(d) progressive excitation of states with flipped spin of an electron and one flipped local spin with a total transfer of angular momentum of  $1 \times \hbar$ . However, the initial state  $|\downarrow; \uparrow \dots \uparrow\rangle$  maintains its probability close to 1, and other states with flipped electron spin and one flipped local spin have much smaller and nonuniform probability (note that increased probability to find a  $|\uparrow; \uparrow \uparrow \uparrow \downarrow\rangle$  state is a consequence of the XXZ chain being open, which brings this state and the  $|\uparrow; \downarrow \uparrow \uparrow \uparrow\rangle$  state into resonance [28,29]). Such peculiar quantum superposition of separable many-body states, with a large contribution from the initial state, leads to local spins maintaining their direction along the  $z$  axis in Fig. 3(c). This can be contrasted with naive (i.e., not taking into account superpositions) intuition [22,35] where a spin- $\downarrow$  electron simply flips the first local spin—the flip then propagates to displace transversally other local spins away from the anisotropy axis, eventually exciting the white spectrum [22] of lowest-energy magnons [28,29].

The same effects—entanglement of an electron state and a quantum state of all local spins [Fig. 4(b)], and thereby in-

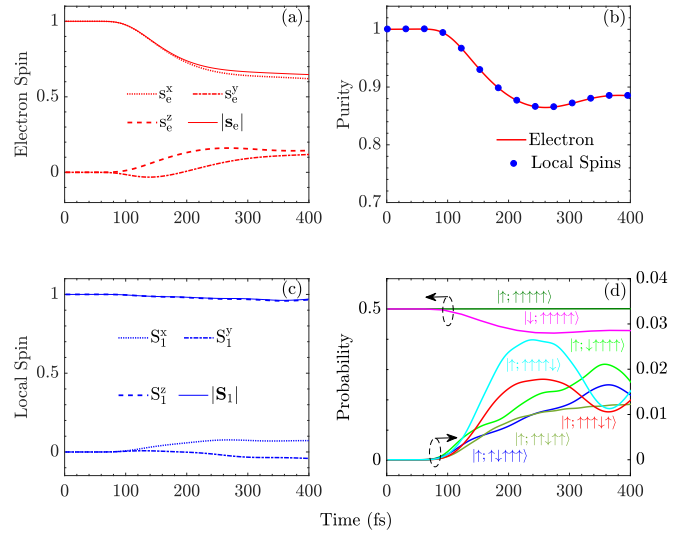


FIG. 4. Panels (a)–(d) plot the same information as panels (a)–(d), respectively, in Fig. 3 but for an injected electron wave packet that at  $t = 0$  is spin-polarized in the  $+x$  direction, i.e., *noncollinear* to local spins pointing in the  $+z$  direction.

duced true decoherence [32,33] of an electron spin [Fig. 4(a)] and local spins [Fig. 4(c)], and a high probability [Fig. 4(d)] of finding the initial state of an electron-spin+local-spins subsystem over the course of time evolution—are present also in the standard STT geometry with noncollinearity between spins of the injected electron and local spins. However, in the standard STT geometry we also find the usual magnetization rotation, i.e.,  $S_1^x \neq 0$  and  $S_1^y \neq 0$  in Fig. 4(c). The probabilities  $\text{prob}_{\text{spins}}^{\text{e+local}}(t)$  in Fig. 4(d) to excite states of the type  $|\uparrow; \uparrow \dots \downarrow \dots \uparrow\rangle$  are simply half of those obtained for collinear geometry in Fig. 3(d) since spin of the injected electron along the  $+x$  direction used in Fig. 4 means  $|\rightarrow_e\rangle = \frac{1}{\sqrt{2}}(|\uparrow_e\rangle + |\downarrow_e\rangle)$ , where only the  $\frac{1}{\sqrt{2}}|\downarrow_e\rangle$  term, entering as a factor of the initial separable many-body state  $|\Psi(t=0)\rangle \equiv |G\rangle \otimes \frac{1}{\sqrt{2}}(|\uparrow_e\rangle + |\downarrow_e\rangle) \otimes |\uparrow_1 \dots \uparrow_N\rangle$ , induces time evolution of local spins and transfer of angular momentum. On the other hand, the  $\frac{1}{\sqrt{2}}|\uparrow_e\rangle \otimes |\uparrow_1 \dots \uparrow_N\rangle$  term in the initial many-body state is an eigenstate [Fig. 2(c)] of an electron-spin+local-spins subsystem, and therefore it has time-independent  $\text{prob}_{\text{spins}}^{\text{e+local}}(t) = 1/2$  in Fig. 4(d). Thus, the identical profile of curves in Figs. 3(d) and 4(d) reveals that in a fully quantum many-body picture there is no difference between the standard STT and quantum STT—both originate from the  $|\downarrow_e\rangle$  factor state brought into the initial many-body state by either minority-spin electrons [35] in spin valves with parallel magnetizations [16], or by the  $|\downarrow_e\rangle$  term in the quantum superposition of electron spin states generated by the polarizing effect of the fixed magnetization that is noncollinear to free magnetization.

We note that increasing the total number of local spins from  $N = 5$  considered here to realistically large values does not change these findings due to Schmidt decomposition [36], which makes it possible to rewrite *any* entangled state of spin- $\frac{1}{2}$  as the sum of just two terms,  $a|\nearrow\rangle \otimes |\Sigma_1\rangle +$

$b|\nearrow\rangle \otimes |\Sigma_2\rangle$ , where  $|\Sigma_{1,2}\rangle \in \mathcal{H}_{\text{spin}}^1 \cdots \otimes \mathcal{H}_{\text{spin}}^N$ . Thus, the Schmidt decomposition state looks the same as the entangled state of spin- $\frac{1}{2}$  with just a single macrospin. We provide counterparts of Figs. 3 and 4 in the Appendix where a single local spin- $\frac{5}{2}$  is employed as the macrospin.

We emphasize that a piecewise-linear increase of differential resistance with the bias voltage observed in Ref. [16] and our explanation of its origin cannot be accounted for by inelastic electron-magnon scattering [37] often observed in magnetic tunnel junctions, which actually leads to the opposite effect in which differential resistance of the junction with collinear magnetizations (parallel or antiparallel) decreases [38,39] with the bias voltage due to opening of additional conduction channels [37].

Our model in Fig. 1 can be viewed as an addition to the atlas of toy models considered in Ref. [26] to explain conventional STT of Slonczewski [2] and Berger [3] due to a *single* injected spin-polarized electron in a plane-wave orbital state that is treated quantum-mechanically, while the magnetization receiving the torque is treated as a classical vector of fixed length. In contrast, in our toy model of Fig. 1 *both* the single injected spin-polarized electron in the wave-packet orbital state and localized spins receiving the torque are treated as a single quantum many-body system. However, to explain conventional STT in realistic junctions requires us to also sum over incoming momenta of all Fermi surface electrons [26,27], which reduces the transverse component of the transmitted and reflected spin currents to nearly zero due to substantial phase cancellation [26,27], so that conventional STT is very nearly proportional to the transverse piece of the incident spin current. While such an effect is not considered in our strictly one-dimensional system (in three-dimensional spin valves, one sums different incident transverse wave vectors that are parallel to junction interfaces [26,27]), we do effectively include averaging over different momenta of plane waves comprising the wave packet. We relegate to future studies an investigation of the nonclassical state of magnetization interacting with a flux of electrons with different momenta arising from different parts of the Fermi surface.

Finally, during the completion of this work we became aware of two studies [40,41] in which magnetization dynamics in collinear spin valves at cryogenic temperatures is treated quantum-mechanically to find quantum STT on a single macrospin signified by  $S_z$  shrinking in length while  $S_x = S_y = 0$ . However, these studies do not invoke the entanglement-induced true decoherence mechanism we discuss (for the case of macrospin in Appendix).

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#### APPENDIX: QUANTUM DESCRIPTION OF SPIN TRANSFER TORQUE AND THE ENSUING DYNAMICS OF A SINGLE MACROSPIN

In this Appendix, we present Figs. 5 and 6, which are the counterparts of Figs. 3 and 4, respectively, for an injected

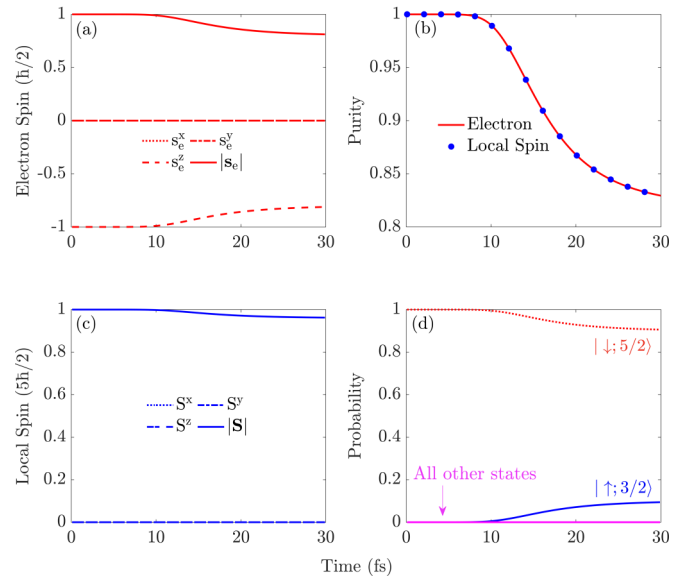


FIG. 5. Time dependence of the expectation value of spin (in units  $\hbar/2$ ) obtained from respective density matrices for (a) spin of injected electron wave packet that at  $t = 0$  points in the  $-z$  direction that is *collinear* and antiparallel to a single local spin- $\frac{5}{2}$  pointing in the  $+z$  direction; and (c) local spin- $\frac{5}{2}$ . (b) Purity defined in Eq. (7) for the subsystem composed of electron degrees of freedom (orbital and spin) or of the subsystem composed of all local spins. (d) Probability in Eq. (8) to find an electron-spin+local-spin subsystem in a many-body quantum state  $|\sigma_e; S_z\rangle$ .

spin-polarized electron wave packet whose spin interacts via  $s$ - $d$  exchange interaction with a macrospin representing the active ferromagnetic layer that receives STT. Note that this Appendix can also be viewed as a rigorous analysis of heuristic arguments provided in Sec. III B of the supplemental material of Ref. [16].

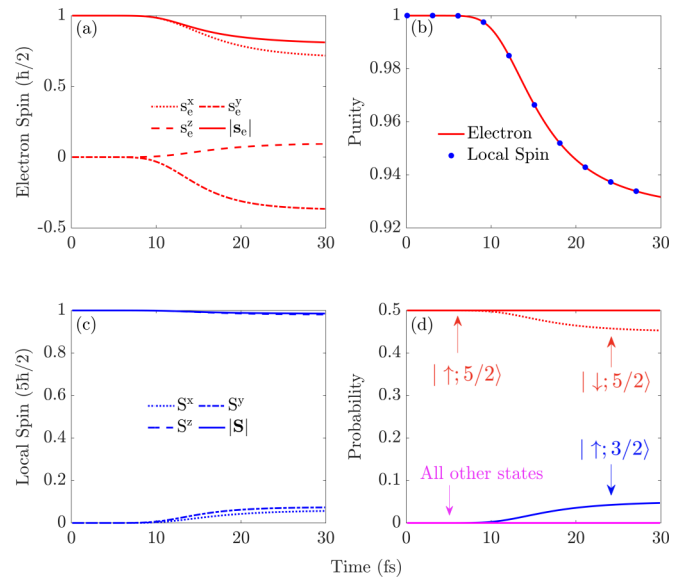


FIG. 6. Panels (a)–(d) plot the same information as panels (a)–(d), respectively, in Fig. 5 but for an injected electron wave packet that at  $t = 0$  is spin-polarized in the  $+x$  direction, i.e., *noncollinear* to local spin- $\frac{5}{2}$  pointing in the  $+z$  direction.

The macrospin is modeled using a single (i.e.,  $N = 1$  in Fig. 1) local spin- $\frac{5}{2}$ . Since we assume an absence of spin-orbit coupling, the orbital state of the wave packet,  $|G\rangle$ , and the state of the electron-spin+local-spins subsystem,  $|\text{spins}\rangle$ , remain uncorrelated [31] over the course of time evolution,  $|\Psi(t)\rangle = |G\rangle \otimes |\text{spins}\rangle$ .

In the case of a  $|G\rangle \otimes |\uparrow\rangle$  injected electron wave packet, with electron spin pointing in the  $+z$  direction, and local spin- $\frac{5}{2}$  being collinear to it in the state  $|S_z = 5/2\rangle$ , there is no STT or entanglement since  $|\Psi(t)\rangle = |G\rangle \otimes |\uparrow\rangle \otimes |5/2\rangle$  is a separable eigenstate of the many-body Hamiltonian, and therefore it does not change with time.

In the case of a  $|G\rangle \otimes |\downarrow\rangle$  injected electron wave packet, with electron spin pointing in the  $-z$  direction and local spin- $\frac{5}{2}$  being collinear in state  $|S_z = 5/2\rangle$ , initial and final states are given by

$$|\Psi(t = 0)\rangle = |G\rangle \otimes |\downarrow\rangle \otimes |5/2\rangle, \quad (\text{A1a})$$

$$|\Psi(t > 0)\rangle = a(t)|G\rangle \otimes |\downarrow\rangle \otimes |5/2\rangle + b(t)|G\rangle \otimes |\uparrow\rangle \otimes |3/2\rangle, \quad (\text{A1b})$$

where the time evolution of probabilities in Fig. 5(d) is given by  $|a(t)|^2$  (red dotted line) and  $|b(t)|^2$  (blue solid line). The total spin in the  $z$  direction remains conserved during time evolution. Over the course of time evolution, local spin- $\frac{5}{2}$  in Fig. 5(c) *does not* rotate away from the  $z$  axis, but it is

shrinking due to true decoherence caused by entanglement in Eq. (A1b).

In the case of a  $|G\rangle \otimes \frac{1}{\sqrt{2}}(|\uparrow\rangle + |\downarrow\rangle)$  injected electron wave packet, with electron spin pointing in the  $+x$  direction and local spin- $\frac{5}{2}$  pointing in the  $+z$  direction in the state  $|S_z = 5/2\rangle$ , initial and final states are given by

$$|\Psi(t = 0)\rangle = |G\rangle \otimes \frac{1}{\sqrt{2}}(|\uparrow\rangle + |\downarrow\rangle) \otimes |5/2\rangle, \quad (\text{A2a})$$

$$|\Psi(t > 0)\rangle = \frac{1}{\sqrt{2}}|G\rangle \otimes |\uparrow\rangle \otimes |5/2\rangle + \frac{a(t)}{\sqrt{2}}|G\rangle \otimes |\downarrow\rangle \otimes |5/2\rangle + \frac{b(t)}{\sqrt{2}}|G\rangle \otimes |\uparrow\rangle \otimes |3/2\rangle, \quad (\text{A2b})$$

where the time evolution of probabilities in Fig. 6(d) is given by  $1/2$  (red solid line),  $|a(t)|^2/2$  (red dotted line), and  $|b(t)|^2/2$  (blue solid line). The latter two are just half of the probabilities in Fig. 5(d). Over the course of time evolution, this change leads to local spin- $\frac{5}{2}$  in Fig. 6(c) *both* rotating away from the  $z$  axis, as in the standard STT of Slonczewski [2] and Berger [3], and shrinking due to true decoherence caused by entanglement in Eq. (A2b). The latter effect is the same as the only consequence of quantum STT in Fig. 5(c).

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