Nonlocal damping of helimagnets in one-dimensional interacting electron systems

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We investigate magnetization relaxation in interacting one-dimensional helimagnetic systems. Relaxation results from the emission of plasmonic excitations into the itinerant electron system due to slow changes of the magnetization profile. This dissipation mechanism leads to a highly nonlocal form of magnetization damping that is strongly dependent on the electron-electron interaction. Forward scattering processes lead to a spatially constant damping, whereas backscattering processes produce a spatially oscillating damping. Due to the nonlocal damping, the thermal fluctuations become spatially correlated over the entire system. We calculate the characteristic magnetization relaxation times for chains of magnetic Fe atoms, magnetic quantum wires, and nuclear helimagnets.

Recently, intense interest has developed in the helical magnetic ordering of one-dimensional (1D) systems of local moments coupled to itinerant electrons. Such systems exhibit a variety of intriguing many-body phenomena, such as spin-Peierls instabilities and induced topological superconductivity, which result from the interplay between magnetism, electronic structure, and interactions. These phenomena may be relevant for a wide variety of physical systems, ranging from magnetic atoms on superconducting and normal metal substrates, to single walled carbon nanotubes and semiconductor-based quantum wires.

While much of the work in this area so far has focused on static and thermodynamic properties of the 1D helimagnets, a richer understanding may be gained by developing and employing new dynamical probes for assessing the behaviors of these systems. For example, an interesting self-tuning effect was proposed for systems dominated by a Ruderman-Kittel-Kasuya-Yoshida (RKKY)-type interaction: the local moments are predicted to order into a spiral arrangement which, through coherent backscattering, gaps out one spin channel of the itinerant electron system for any value of the electron density. This remarkable phenomenon was even suggested as providing a route towards realizing topologically protected Majorana bound states in quantum wires. However, because direct probes of magnetization are unavailable for many systems, it can be challenging to positively identify this intriguing magnetic state necessarily via indirect means.

With further theoretical understanding of dynamical responses, such as typical damping or relaxation times, additional tests (e.g., density quenches which change the preferred ordering wave vector) could be used to clarify the natures of the underlying states.

More generally, magnetization relaxation processes determine the magnetic response to external perturbations as well as to spontaneous thermal fluctuations. Furthermore, the nature of the magnetic response is crucially important for noise and magnetization dynamics in magneto-electronic devices. A better understanding of the spin dynamics in 1D helimagnets may pave the way for exploring phenomena such as current-driven magnetization dynamics, with potential practical applications beyond those envisaged so far. Thus, the investigation of microscopic damping mechanisms is essential for developing a thorough fundamental and practical understanding of these exciting new magnetic systems.

Given the motivations above, in this work we investigate the relaxation of 1D helimagnets via the emission of plasmon excitations into the interacting itinerant electron system. Interestingly, previous theoretical works have predicted that electron-electron interactions in such 1D systems may play important roles both in establishing ordering and in the relaxation dynamics of weakly-coupled (non-ordered) nuclear spins. In this work, we use a bosonization approach to study the non-perturbative effects of electron-electron interactions on the damping of ordered spins. We find that interactions have a profound effect on damping, leading to an enhancement of the damping by several orders of magnitude. The damping has a highly non-local character. Consequently, the thermal fluctuations become spatially correlated over the entire sample. We calculate the characteristic magnetization relaxation times caused by the plasmon excitations for three classes of systems: chains of magnetic Fe atoms, (Ga,Mn)As quantum wires, and nuclear helimagnets formed in InAs quantum wires.

We now calculate the magnetization-damping rate, relating it to the energy absorption rate of the itinerant electron system subjected to small fluctuations about a static helimagnetic profile. Our approach is based on...
the theoretical framework developed for magnetization damping in metallic ferromagnets [15]. A key ingredient of the model is that the dynamics of the low-lying collective spin excitations (the Goldstone modes) are parametrized by a classical magnetization order-parameter field whose magnitude is assumed to be constant in time and homogeneous in space, while its local orientation is allowed to fluctuate. In this case, the evolution of the spin system can be described by the Landau-Lifshitz-Gilbert (LLG) phenomenology [14, 15, 21].

In describing the interaction between the magnetic order parameter field and the itinerant electron system, we assume that the magnetization evolves very little over the characteristic timescales of electron dynamics. In this case, the response of the electron system can be calculated using linear response (Kubo) theory.

To begin, we parametrize the local spin-order by the unit vector \( \mathbf{m}(z,t) \) oriented parallel to the magnetization vector: \( \mathbf{M}(z,t) = M_s \mathbf{m}(z,t) \). Throughout this work, we use coordinate axes with the z-axis oriented along the 1D conductor (see Fig. 1). The evolution of the magnetization is described by the LLG equation [14, 15]:

\[
\dot{\mathbf{m}}(z,t) = -\gamma \mathbf{m}(z,t) \times [\mathbf{H}_{\text{eff}}(z,t) + \mathbf{h}_T(z,t)] + \mathbf{m}(z,t) \times \int dz' \alpha(z,z') \dot{\mathbf{m}}(z',t),
\]

where \( \dot{\mathbf{m}} = \partial \mathbf{m}/\partial t \), \( \gamma = g \mu_B / h \) is the gyromagnetic ratio in terms of the g-factor of local spins and the Bohr magneton \( \mu_B \), and \( \mathbf{H}_{\text{eff}} = -\delta F/\delta \mathbf{M} \) is the effective field found by varying the magnetic free energy functional \( F[\mathbf{M}] \) with respect to the magnetization. The quantity \( \mathbf{h}_T(z,t) \) in the first line is a stochastic magnetic field induced by the thermal fluctuations (to be discussed further below). We assume that the free energy functional stabilizes an equilibrium helimagnetic texture of the form [5, 8, 10–12]

\[
\mathbf{m}_0(z) = [\cos(qz), \sin(qz), 0],
\]

where \( q \) depends on the ordering mechanism.

Magnetoelastic relaxation is described by the second-rank Gilbert damping tensor \( \alpha_{ij}(z,z') \) in Eq. (1). We relate the Gilbert damping to the total energy dissipation \( \dot{E} = \int dz \mathbf{M} \cdot \delta F/\delta \mathbf{M} \) using Eq. (1), which gives

\[
\dot{E}(t) = -\frac{M_s}{\gamma} \int dz dz' \dot{\mathbf{m}}(z,t) \cdot [\dot{\alpha}(z,z') \mathbf{m}(z',t)].
\]

Due to energy conservation, the energy lost by the magnetic system must be gained by the itinerant electron system to which it is coupled. This implies that \( \dot{E} = -(H(t)) \), where \( H(t) \) is the Hamiltonian of the itinerant electron system coupled to the magnetization \( \mathbf{m}(z,t) = \mathbf{m}_0(z) + \delta \mathbf{m}(z,t) \):

\[
H = \int dz \psi^\dagger(z) \left[ \frac{\hbar^2}{2m} + h_0 \mathbf{m}(z,t) \cdot \mathbf{\sigma} \right] \psi(z) + \frac{1}{2} \int dz dz' \psi^\dagger_{\sigma}(z) \psi_{\sigma}(z') V_{\text{ee}}(z-z') \psi_{\sigma}(z') \psi_{\sigma}(z).
\]

Here, \( \psi(z) = [\psi_T(z), \psi_L(z)] \) is the spinor-valued fermionic field operator, \( \tilde{\mathbf{p}}_z \) is the momentum operator, \( \psi_{\text{ee}} \) is the electron-electron interaction potential, \( h_0 \) is the magnetic coupling, and \( \mathbf{\sigma} \) is the vector of Pauli matrices. Here and below, summation over repeated indices is implied.

We now use linear response theory to obtain the rate of change of the electronic energy due to a slow evolution of the magnetization \( \mathbf{m}(z,t) \):

\[
\langle \dot{H}(t) \rangle = -\frac{i}{\hbar} \int_{-\infty}^{\infty} dt' \theta(t-t') \langle [\dot{H}(t), \delta H(t')] \rangle,
\]

where \( \theta(t) \) is the Heaviside step function, and \( \delta H(t') = (2h_0/\hbar) \int dz \dot{\mathbf{m}}(z,t) \cdot \dot{\mathbf{s}}(z,t) \) is the perturbing Hamiltonian produced by the small variation \( \delta \mathbf{m}(z,t) \) of the helimagnetic order. Here, \( \dot{\mathbf{s}}(z,t) \) is the spin-density operator \( \dot{\mathbf{s}}(z) = (\hbar/2) \psi^\dagger(z) \mathbf{\sigma} \psi(z) \), taken in the interaction picture with respect to the unperturbed Hamiltonian.

Fourier transforming Eq. (5) with respect to time and using \( \dot{H}(t) = (2h_0/\hbar) \int dz \dot{\mathbf{m}}(z,t) \cdot \dot{\mathbf{s}}(z,t) \) yields

\[
-i\omega \langle [\dot{\mathbf{H}}(\omega), \delta \mathbf{H}(0)] \rangle = (2h_0/\hbar)^2 (1/2\pi) \int dz dz' \int d\omega' i(\omega - \omega') \langle \mathbf{\chi}_{ij}(z,z',\omega')/\omega' \rangle \langle \delta m_j(z',\omega') \rangle \langle \delta m_j(z',\omega') \rangle,
\]

where \( \mathbf{\chi}_{ij}(z,z',\omega) = \int dt \chi_{ij}(z,z',t) \exp(i\omega t) \) is the Fourier transform of the spin susceptibility \( \chi_{ij}(z,z',t) = -(i/\hbar) \theta(t) [\mathbf{s}_i(z,t), \mathbf{s}_j(z',0)] \). To leading order in the precession frequency, the behavior of the energy change is captured by replacing \( i\chi_{ij}(z,z',\omega) \) by its value in the zero frequency limit. Transforming back to the time domain and using \( \delta \dot{m}_j = \dot{m}_j \) gives the energy absorption rate

\[
\langle \dot{H}(t) \rangle = \frac{4h_0^2}{\hbar^2} \int dz dz' \dot{m}_j \left[ \lim_{\omega \rightarrow 0} \frac{i\chi_{ij}(z,z',\omega)}{\omega} \right] \dot{m}_j.
\]

Comparing Eq. (6) with Eq. (4), we identify the following expression for the Gilbert damping tensor [15]:

\[
\alpha_{ij}(z,z') = -\frac{4\gamma h_0^2}{\hbar^2 M_s} \lim_{\omega \rightarrow 0} \frac{3m \chi_{ij}(z,z',\omega)}{\omega}.
\]

Next, we explicitly calculate the Gilbert damping tensor in Eq. (7). To facilitate calculation of the spin susceptibility, we transform to a non-uniformly rotated frame via the unitary transformation \( \psi = U(z) \psi' \), with

\[
U(z) = e^{i\varphi z / \sqrt{2}}.
\]

This transformation "untwists" the helix, rendering the free electron part of the transformed Hamiltonian \( U_H = U H U^\dagger \) translationally invariant,

\[
H^{(0)} = \int dz \psi_{\text{ee}}^\dagger(z) \left[ \frac{\hbar^2}{2m} + \frac{h_0}{2m} \mathbf{\sigma} \tilde{\mathbf{p}}_z + h_0 \sigma_z \right] \psi_{\text{ee}}(z),
\]

while the interaction term is unaffected. In this representation, the spin susceptibility and the Gilbert damping tensor transform to \( \chi^{(0)}(z,z',t) = R(z) \chi(z,z',t) R^\dagger(z') \) and \( \alpha^{(0)}(z,z',t) = R(z) \alpha(z,z',t) R^\dagger(z') \), where \( R(z) \) is the SO(3) matrix associated with \( U(z) \). The energy dispersion of \( H^{(0)} \) is shown in Fig. 2. Its eigenfunctions are \( \psi_{n,k}(z) = \eta_{n,k} \otimes \psi_{n,k}(z) \), where \( n \in \{1, 2\} \) is the band index, \( \eta_{n,k} \) is the eigenspinor, and \( \psi_{n,k}(z) = \exp(ikz)/\sqrt{L} \), for a system of length \( L \).
In this work, we set the chemical potential in the gap between bands around \( k = 0 \) such that the single-particle dispersion in Eq. (3) features only a single branch of right and left moving modes at the Fermi energy. We neglect interband couplings and write an effective description within the lowest band, linearized about the Fermi wavevectors \( \pm k_F \). We fix the spinor parts of the wave functions to their values at the Fermi energy (Fig. 2).

To compute the spin-spin susceptibility in the presence of electron-electron interactions, we employ a bosonic description. As a first step, we express the fermionic field operator (projected into the lowest band) as a superposition of fields representing right (+) and left (-) movers:
\[ \psi_u(z) = \psi_+(z) + \psi_-(z) \]
The fields \( \psi_r \) (\( r \in \{+, -\} \)) take the form \( \psi_r = \eta_r \otimes \psi_r(z) \), where \( \eta_r = \eta_{k,r} \) and the spatial part (in terms of the destruction operators \( \eta_{k,r} \)) is given by:
\[ \psi_r(z) = \frac{1}{\sqrt{\pi k_F a}} \int k_F e^{ik_F z} \sum_k \eta_{k,r} e^{-ik_F z} \]
Substituting the fermionic field operator into the Hamiltonian \( H_u \), performing a Fourier transformation to k-space, and evaluating \( V_{ee}(q) \) at momentum zero and \( 2k_F \) for forward- and back-scattering processes, respectively, we obtain
\[ H_u = \sum_{k,r} v_F k_F \eta_{k,r} e^{ik_F z} + \sum_{q,r} \left( g_2 \rho_{q,r} \rho_{q,-r} + g_4 \rho_{q,r} \rho_{-q,-r} \right) \]
Here, \( v_F = \hbar k_F / m - \hbar q \eta_{k,r} / 2m \), \( g_2 = (V_{ee}(0) - |\eta_{k,r}|^2 V_{ee}(k_F))/2L \), \( g_4 = V_{ee}(0)/2L \), and \( \rho_{q,r} = \sum_k c_{-q,r}^\dagger c_{k,r} e^{ik_F z} \) is the Fourier-transformed density operator for right/left-movers. Following the standard procedure [23], we write Eq. (9) in the bosonized form
\[ H_u = \frac{\hbar}{2\pi} \int dz \left[ u_{\text{eff}} (\partial_z \phi)^2 + \frac{u_{\text{eff}}}{K} (\partial_z \phi)^2 \right] , \]
where \( \phi \) and \( \theta \) are the bosonic fields, \( u_{\text{eff}} \) is the density wave velocity, and \( K \) is the Luttinger parameter.

The bosonic representations of the fermionic fields are
\[ \psi_r(z) = \eta_r \otimes \frac{U_r}{\sqrt{2\pi a}} e^{i k_F z} e^{-i \phi(z) - i \theta(z)} , \]
where \( a \) is an infinitesimal short distance cutoff [24] and \( \{U_r\} \) are the Klein factors. The repulsive electron-electron interaction implies that \( 0 < K \leq 1 \), where \( K = 1 \) for non-interacting electrons.

We calculate the spin susceptibility following the standard approach for Luttinger liquids, see Ref. [23] for details. The resulting (imaginary) time-ordered spin-spin correlation function [25], \( \chi_{u,ii}(z, z', \tau) = -i \langle T_{\tau} \delta_{\text{ex}}(z, \tau) \delta_{\text{ex}}(z', 0) \rangle \), is diagonal and can be written as
\[ \chi_{u,ii}(z, \tau) = \chi_{u,ii}^{(0)}(z, \tau) + \chi_{u,ii}^{(2k_F)}(z, \tau) \cos(2k_F z), \]
where
\[ \chi_{u,ii}^{(0)} = - \left( \frac{\hbar}{2\pi} \right)^2 K (1) \Lambda_{ii}^{++} / 2 \left( v^2 - \frac{\bar{u}}{2} \right)^2, \]
\[ \chi_{u,ii}^{(2k_F)} = - \left( \frac{\hbar}{2\pi} \right)^2 \frac{\Lambda_{ii}^{++}}{2a^2} \left( \frac{\bar{a}^2 + \bar{u}^2}{2} \right)^2 K \]
Here, \( \bar{u} = u_{\text{eff}} \pi / L \), and \( \Lambda_{ii}^{++} = |\eta_{k,r}|^2 \). The spin susceptibility is \( \chi_{u,ii}(z, \tau) = -(2/\hbar) \delta(\tau) \sum_{l} \chi_{u,ii}(z, \tau) \), where \( \chi_{u,ii}(z, \tau) \) is the time-ordered correlation function in real time, which is obtained via the Wick rotation \( \tau = it + \theta \). For \( K < 1 \), \( \chi_{u,ii}^{(2k_F)}(\omega) / \omega \) diverges in the low-frequency limit. We regularize the divergence by evaluating the expression at the low-frequency cut-off \( \omega_0 = u_{\text{eff}} \pi / L \) set by the finite length of the system [26].

The analysis above gives the Gilbert damping tensor
\[ \alpha_{u,ii}(z, z') = \alpha_{u,ii}^{(0)} + \chi_{u,ii}^{(2k_F)}(2k_F z), \]
\[ \alpha_{u,ii}^{(0)} = \gamma \frac{K (1) \Lambda_{ii}^{++}}{2 \pi \hbar M_u u_{\text{eff}}}, \]
\[ \alpha_{u,ii}^{(2k_F)} = \gamma \frac{K (1) \Lambda_{ii}^{++} - \epsilon K_{-1/2} F_K(\zeta)}{2^{1/2 + K} \pi^{3/2} \hbar M_u u_{\text{eff}}}, \]
Here, \( \Gamma(K) \) is the gamma function, \( \zeta \equiv \omega_0 / \hbar M_u \), and \( F_K(\zeta) = \pi (I_{-1/2}(\zeta) - L_{-1/2}(\zeta)) - 2 \zeta \cos(\pi K) \zeta K_{-1/2}(\zeta) \). The 

Equations [14] - [16] are the central results of this work, and describe magnetization relaxation of interacting 1D helimagnets. The damping consists of two parts with very distinct position dependencies: a homogeneous term \( \sim \alpha^{(0)} \) and a rapidly oscillating term \( \sim \alpha^{(2k_F)} \). The corresponding highly nonlinear magnetization relaxation caused by the plasmon excitations differs markedly from the damping of conventional metallic ferromagnets, which is believed to be local [28].

Constraints of the model reduce the number of independent tensor elements. First, \( \Lambda_{ii}^{++} = \Lambda_{ii}^{-+} = 0 \) implies that the damping tensor is described by four independent coefficients: \( \alpha_{u,xx}^{(0)}, \alpha_{u,zz}^{(0)}, \alpha_{u,xx}^{(2k_F)}, \) and \( \alpha_{u,yy}^{(2k_F)} \). Second, the
constraint \( \mathbf{m} \cdot \mathbf{m} = 0 \) imposed by normalization implies that \( \dot{n}_0 = 0 \) in the rotated reference frame (for small \( \delta \mathbf{m} \)). Thus, the damping is governed by only two coefficients: \( \alpha_{u,zz}^{(0)} \) and \( \alpha_{u,yy}^{(2k_F)} \).

What are the characteristics of these two independent damping coefficients? The tensor element \( \alpha_{u,zz}^{(0)} \) originates from forward scattering processes and is proportional to the momentum-momentum correlator \( \langle \partial_x \theta(\tilde{z}, \tau) \partial_y \theta(0,0) \rangle \), which is inversely proportional to the Luttinger parameter \( K \). Consequently, electron-electron interaction enhances \( \alpha_{u,zz} \) as \( K^{-1} \). A much more complex dependency of the electron-electron interactions is seen in the magnetization damping caused by backscattering processes. Remarkably, electron-electron interactions increase \( \alpha_{u,yy}^{(2k_F)} \) by nearly four orders of magnitude compared to its value in the non-interacting limit \( K = 1 \) (Fig. 3a). The tensor element \( \alpha_{yy}^{(0)} \) reaches a maximum at \( K \approx 0.1 \) before it drops quickly to zero in the strongly interacting regime \( K \rightarrow 0 \). In this limit, the potential energy \( (u_{\text{eff}}/K)(\partial_y \phi)^2 \) of the bosonic Hamiltonian completely governs the electron dynamics and density variations of the Luttinger liquid become insusceptible to time variations in the magnetization. The dramatic enhancement of \( \alpha_{yy}^{(2k_F)} \) is related to that electron-electron interactions make the damping extremely sensitive to \( \zeta \sim a/L \), which measures the ratio of the short distance cut-off to the long distance cut-off (Fig. 3b). In absence of interactions, \( \alpha_{yy}^{(2k_F)} \) approaches a constant value in the limit \( \zeta \rightarrow 0 \). However, with interactions, \( \alpha_{yy}^{(2k_F)} \) is singular in this limit and the singularity becomes stronger with increasing strength of the interactions.

To investigate the experimental consequences of Eqs. (14)-(16), we discuss thermal fluctuations and estimate the characteristic relaxation time for three classes of systems which are proposed to hold 1D helimagnetic states. The magnetic order is assumed to be to stabilized by the RKKY interaction, which implies that \( q = k_F \) [27].

For a magnetization precessing at frequency \( \omega \), Eq. (1) yields two characteristic relaxation times \( \tau_r^{(0.2k_F)} \) associated with \( \alpha^{(0.2k_F)} \) on the order of \( \tau_r^{(0.2k_F)} \sim (\alpha^{(0.2k_F)}L\omega)^{-1} \). Table I shows \( \tau_r^{(0.2k_F)} \) for three systems: chain of magnetic Fe atoms, \( \text{G}_0.98\text{Mn}_{0.02}\text{As} \) quantum wire, and InAs wire in which the nuclear spins are hyperfine coupled to the itinerant system. The (Ga,Mn)As and InAs wire cross sections are assumed to contain 50 × 50 unit cells. The characteristic magnon frequency is \( \omega = k_BT_r/hI_{\perp}^{1/(3-2K)} \) [12], where \( T_r \) is the critical temperature and \( hI_{\perp} \) is the total spin of the cross section. Both the magnetic wire and the magnetic chain have relaxation times in the pico- to nano- second range. In contrast, the relaxation times of the nuclear system are more than a millisecond. This ultra-slow relaxation rate opens the possibility for a direct experimental detection of the relaxation process in nuclear helimagnetic systems.

The form of the damping tensor has remarkable implications for the statistical properties of thermal fluctuations. While the average of the stochastic magnetic field \( \mathbf{h}_T \) in Eq. (1) is zero, \( \langle \mathbf{h}_T \rangle = 0 \), its correlations (in accordance with the fluctuation-dissipation theorem) are given in the classical (Maxwell-Boltzmann) limit as [18, 30]

\[
\langle h_{T,i}(z,t)h_{T,j}(z',t') \rangle = \frac{2k_BT}{\gamma M_s} \alpha_{ij}(z,z') \delta(t-t'), \quad (17)
\]

where \( T \) is the temperature and the average \( \langle ... \rangle \) is taken over an ensemble in thermal equilibrium. According to Eq. (17), the thermal fluctuations are correlated over the entire sample. The fluctuations divide into two distinct classes; one class characterized by a spatially constant

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### Table I: Material parameters (adapted from Ref. [5]) and characteristic relaxation times for three classes of systems.

| Material         | Chain | Magnetic wire | Nuclear wire |
|------------------|-------|---------------|--------------|
| L (\( \mu m \))  | 0.05  | 20            | 20           |
| \( \mu_F \) (meV) | 10    | 20            | 1.0          |
| a (\( \text{nm} \)) | 2     | 5             | 37           |
| K                | 0.5   | 0.5           | 0.5          |
| \( u_{\text{eff}} \) (ms\(^{-1} \)) | 5.9 \times 10^4 | 3.2 \times 10^5 | 1.1 \times 10^7 |
| \( M_s/\gamma \) (Jsm\(^{-1} \)) | 7.0 \times 10\(^{-25} \) | 2.3 \times 10\(^{-23} \) | 5.2 \times 10\(^{-23} \) |
| \( h_0 \) (meV)  | 6     | 5             | 0.2          |
| \( I_\perp \)     | 2     | 500           | 300          |
| \( T_c \) (K)     | 14    | 2             | 0.0074       |
| \( \tau^{(0)} \) (s) | 1.8 \times 10^{-11} | 2.3 \times 10^{-8} | 2.0 \times 10^{-1} |
| \( \tau^{(2k_F)} \) (s) | 1.0 \times 10^{-11} | 7.6 \times 10^{-11} | 4.8 \times 10^{-3} |

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FIG. 3: (Color online). (a) The dimensionless damping parameter \( \alpha^{(2k_F)}(K,\zeta) = \alpha_{u,yy}^{(2k_F)}(K,\zeta)/\alpha_{u,yy}^{(2k_F)}(1,\zeta) \) as a function of the electron-electron interaction parameter \( K \) with \( \zeta = a\omega_0/u_{\text{eff}} \) fixed at \( \zeta = 3.1 \times 10^{-3} \). (b) The damping parameter as a function of \( \zeta \) for different \( K \).
correlation and a second class characterized by an oscillating $\sim \cos(2k_F z)$ correlation. Strongly correlated thermal fluctuations of this form have not been reported or investigated before in any magnetic system, and a thorough investigation of how the associated stochastic field $h_F(z)$ in Eq. (1) influences the magnetization dynamics should be an interesting task for future studies.

In conclusion, we have developed a theoretical formalism of magnetization dissipation of 1D helimagnets via the emission of plasmon excitations. The damping is found to be highly nonlocal and strongly dependent on the electron-electron interaction, and differs markedly from the damping of conventional metallic ferromagnets.

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[24] The cutoff parameterizes a finite bandwidth. In numerical calculations, we use $a = 1/k_F'$ where $k_F'$ is the Fermi vector in the original (non-rotated) reference frame.
[25] Time-ordered quantities are indicated with tildes.
[26] For the cut-off frequency $\omega_0 = 2\pi u_{ab}/L$ to be considered as an infinitesimal quantity, we need $\omega_0 \tau_s << 1$, where $\tau_s \sim a/u_{ab}$ is the correlation time of the spin susceptibility. Using $a \sim 1/k_F$, this implies that the system length should be much larger than the Fermi wavelength, $L >> \lambda_F$.
[27] For RKKY systems, the wave vector $k_F$ in the rotated frame is twice the lab-frame wave vector $k_F'$, i.e., $k_F = 2k_F'$.
[28] Note that a recent first-principle study of Permalloy reports signatures of nonlocality in the damping [31].
[29] An alternative estimate for InAs, which assumes that $\sqrt{N}$ of the total number $N$ of nuclei contained in the cross section of the wire are fully polarized (each with an average spin of 3), yields the relaxation times $\tau^{(0)} = 2.0 \times 10^{-5}$ s and $\tau^{(2k_F)} = 4.8 \times 10^{-7}$ s.
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