Proposal for a Quantum Magnetic *RC* Circuit

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We propose a setup that is the spin analog of the charge-based quantum *RC* circuit. We define and compute the spin capacitance and the spin resistance of the circuit for both ferromagnetic and antiferromagnetic systems. We find that the antiferromagnetic setup has universal properties, but the ferromagnetic setup does not. We discuss how to use the proposed setup as a quantum source of spin excitations, and put forward two possible experimental realizations, using either ultracold atoms in optical lattices or artificially engineered atomic-spin chains.

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Introduction.—The aim in the field of spintronics in insulating magnets [1–9] is to use purely magnetic collective excitations, such as magnons [10] or spinons [11], to perform logic operations in the absence of charge transport. Thereby, it is possible to circumvent the problem of excess Joule heating that occurs due to the scattering of conduction electrons in more traditional electronic devices, leading to lower energy dissipation in such spintronic devices [12]. Since this excess heating is a limiting factor in the design of electronic devices, spintronics in insulating magnets is considered one of the candidates to become the next computing paradigm. Furthermore, the fact that the elementary excitations in ferromagnetic insulators obey bosonic statistics may offer additional benefits [1].

Several experiments that display the capability to create and detect pure spin currents in magnetic insulators have been performed recently. Creation of a magnon current has been shown to be possible using the spin Hall effect [4], the spin Seebeck effect [5], as well as laser-controlled local temperature gradients [13]; detection of magnon currents has been performed using the inverse spin Hall effect [4,14]. However, analogously to quantum optics where the single-photon source is a major element to encode or manipulate a quantum state [15], or to quantum electronics where an on-demand electron source has been recently realized [16–18], a more controllable way of creating quantum spin excitations may ultimately be desirable.

Besides offering great potential for applications, single-excitation sources are also fascinating from a more fundamental point of view. This is illustrated by the single-electron source, which violates the classical laws of electricity [19,20]. Furthermore, in the linear response regime and at low driving frequency, a single-electron source can be described in terms of a quantum *RC* circuit whose charge relaxation resistance has a universal value in the coherent regime [19,21–27] (this universality is destroyed when one moves into the incoherent regime [28]). Motivated by these considerations, we analyze here a setup that we propose could potentially act as an on-demand coherent source of magnons or spinons and compute the equivalent *RC* parameters of such a circuit from a microscopic model.

By drawing analogy to the charge-based quantum *RC* circuit, we propose that the setup depicted in Fig. 1(a) is equivalent to a "quantum magnetic *RC* circuit." We mean by this that in the displayed setup $M_D(\omega)$, the excess magnetization of the magnetic grain or magnetic dot (see below), is related to the applied magnetic field $B_D(\omega)$ by

$$\frac{M_D(\omega)}{B_D(\omega)} = C_M(1 + i\omega C_M R_M). \tag{1}$$

Here, C_M and R_M are the magnetic resistance and capacitance, respectively, of the equivalent *RC* circuit [see Fig. 1(b)]. We emphasize that our proposed magnon or spinon source is not equivalent to a classical spin battery [29] but operates at the quantum level (i.e., on the level of individual coherent magnons or spinons).

The excess magnetization of the nonitinerant magnetic dot is defined as $M_D(\omega) = g\mu_B N_D(\omega)$, where $N_D(\omega)$ is the Fourier transform of $N_D(t)$ which is the time-dependent excess number of magnetic quasiparticles in the dot. These quasiparticles are the elementary quantum excitations (magnons or spinons) of the Heisenberg Hamiltonian which we will use to describe the dot. We must make a clear distinction between ferromagnetic (FM) and antiferromagnetic (AF) systems here. The main difference between the two lies in the different statistics obeyed by the respective elementary excitations, whereas the FM magnons obey bosonic statistics, the AF spinons in contrast behave according to fermionic statistics. This leads us to expect very different behavior between these systems.

Model.—Our setup consists of a magnetic dot or magnetic grain that is weakly exchange coupled to a large magnetic reservoir. Both the magnetic dot and reservoir are assumed to be nonitinerant magnets, described by a

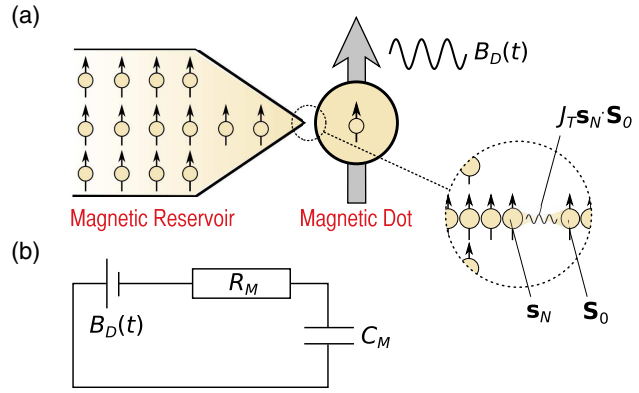

FIG. 1 (color online). (a) Schematic representation of the setup. The weakly coupled reservoir and dot are both modeled as 1D chains in this work. Parallel collections of such 1D chains are realized in bulk materials such as $SrCuO_2$ (Ref. [30]) and Cs_2CoCl_4 (Ref. [31]), as well as in ultracold atoms in optical lattices [32,33]. (b) Equivalent circuit representation of the setup; see Eq. (1) for the definition of R_M and C_M.

Heisenberg Hamiltonian. For concreteness, we will model our subsystems as 1D spin chains. We characterize the system by the Hamiltonian $H = H_0 + H_T$. Here, $H_0 = H_D + H_R$ describes the isolated subsystems and H_T the weak magnetic exchange interaction between the dot and reservoir. The Heisenberg Hamiltonian H_D that describes the isolated magnetic dot is given by

$$H_D = \sum_{\langle ij \rangle} \mathbf{S}_i \cdot \mathbf{J}_D \cdot \mathbf{S}_j + g\mu_B \sum_i [\mathbf{B}_D^0 + \mathbf{B}_D(t)] \cdot \mathbf{S}_i. \quad (2)$$

$\mathbf{J}_D$ denotes a diagonal 3×3 matrix with $\mathrm{diag}(\mathbf{J}_D) = J_D\{1, 1, \Delta_D\}$. J_D is the magnitude of the exchange interaction and Δ_D the anisotropy in our model. $J_D \lesseqgtr 0$ corresponds, respectively, to the FM and the AF ground state. $\mathbf{B}_D^0 = B_D^0 \mathbf{e}_z$ and $\mathbf{B}_D(t) = B_D(t)\mathbf{e}_z$ are, respectively, the static- and time-dependent component of the magnetic field applied to the dot. The Hamiltonian H_R that describes the reservoir is given by Eq. (2) with parameters $\mathbf{J}_R, J_R, \Delta_R$, B_R^0, and $B_R(t) = 0$. We will use a lowercase $\mathbf{s}_i$ to denote the ith spin in the reservoir. H_T will be defined later.

We can either use the Holstein-Primakoff [10] (for FM systems) or the Jordan-Wigner [11] (for spin-1/2 AF systems) transformation to map the spin-ladder operators on respectively bosonic or fermionic creation or annihilation operators, corresponding to spinless quasiparticles with magnetic moment $g\mu_B \mathbf{e}_z$ (see Ref. [34]). Regardless of the statistics of the quasiparticles, we will denote an annihilation operator in the reservoir (dot) by $r_i(d_i)$.

In thermal equilibrium, the ground state of an isolated dot contains a fixed number $N_0 = \sum_i \langle d_i^\dagger d_i \rangle_0$ of magnetic quasiparticles. Here, $\langle \ldots \rangle_0$ denotes the average with respect to the Hamiltonian H_0 with $B_D(t) = 0$. We now define the excess number of magnetic excitations on the dot as $\hat{N}_D = \sum_i d_i^\dagger d_i - N_0$.

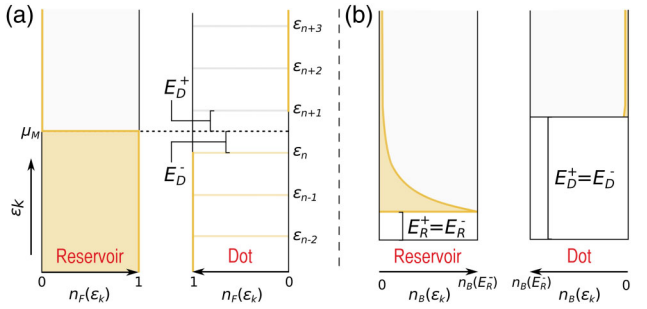

FIG. 2 (color online). (a) Band structure of the AF system, assuming a small AF dot with discrete energy levels. Since the magnetic quasiparticles obey fermionic statistics, all states above (below) μ_M are empty (filled) at $T = 0$. We define $E_D^\pm = \varepsilon_{n+1/n} - \mu_M$. (b) Band structure of the FM system at finite $T \ll E_D^+$. The quasiparticles obey bosonic statistics. Here, $E_{D(R)}^+ = E_{D(R)}^-$ denotes the energy of the lowest energy level in the dot (reservoir).

We will consider magnetic dots whose Hamiltonian can be diagonalized as $H_D = \sum_k (\varepsilon_k - \mu_M) d_k^\dagger d_k$, with ε_k the dispersive energy of the excitations and μ_M the magnetic equivalent of the chemical potential. The parameters R_M, C_M of the quantum RC model are only well defined if adding and removing a quasiparticle from the dot involves a finite amount of energy, i.e., if the spectrum has a gap $E_D^\pm$ that satisfies $|E_D^\pm| \gg k_B T, \hbar\omega, J_T$ (see Fig. 2). In small magnetic dots of size L, quantization of the wave vector k in multiples of $2\pi/L$ leads to a level splitting (and hence $E_D^\pm$; see Fig. 2) of order $J_D(a/L)^2$ for FM dots, and $J_D a/L$ for AF dots (a is the nearest-neighbor distance).

In principle, an anisotropy $\Delta_D > 1$ gives rise to a bulk gap in large AF dots as well. However, the resulting system is the magnetic equivalent of a Mott insulator, rather than the equivalent of a band insulator [11]. As a consequence, the excitations are no longer the $d_k^{(\dagger)}$'s of the original model, and the resulting model does not allow for a straightforward analysis. The opening of a bulk gap $E_D^\pm$ by an applied magnetic field requires a staggered field, with wave vector $2k_F \approx \pi/a$. Since neither mechanism allows us to create a bulk gap $E_D^\pm$ for large AF dots in a straightforward manner, we will rather focus on small AF dots with a finite level splitting due to the quantization of the wave vector k for AF systems. The magnetic chemical potential is given by $\mu_M = g\mu_B |B_D^0|$.

For FM dots, we put $\mu_M = 0$. There exist two mechanisms that allow for a finite gap even in large FM dots. First, an anisotropy $\Delta_D > 1$ gives rise to a gap $E_D^+ = 2|J_D|S_D(\Delta_D - 1)$ [see Fig. 2(b) for the definition of E_D^+ for FM systems]. Second, application of a magnetic field B_D^0 leads to $E_D^+ = g\mu_B |B_D^0|$.

For FM subsystems, we will assume that the reservoir is described by the isotropic Heisenberg Hamiltonian, i.e., with $\Delta_R = 1$. For AF subsystems, we will assume initially that the reservoir as well as the dot are easy-axis AF spin-1/2 spin chains, i.e., with $\Delta_{R(D)} = 0$. This has the

advantage that the excitations can be mapped on free fermions. We will show later how to extend our results for spin chains with finite anisotropy.

The exchange interaction between dot and reservoir is given by $J_T \mathbf{s}_N \cdot \mathbf{S}_0$, where $\mathbf{s}_N$ denotes the last spin in the reservoir, and $\mathbf{S}_0$ the first spin in the dot. J_T is the smallest parameter in the problem, and we analyze the effect of H_T using perturbation theory. We show in the Supplemental Material [34] that the out-of-plane component of the interaction does not significantly affect our results, so that we can approximate H_T by $\hat{H}_T = J_T[r_N^\dagger d_0 + r_N d_0^\dagger]$.

By using linear-response theory, we can calculate the change in magnetization $M_D(\omega)$ due to a small time-dependent change in $B_D(\omega)$. It is given by $M_D(\omega) = (g\mu_B)^2 G_{\mathrm{ret}}(\omega) B_D(\omega)$.

The retarded Green's function is given by $G_{\mathrm{ret}}(\omega) = G(\omega) + G^*(-\omega^*)$, where

$$G(\omega) = i\int_0^\infty \mathrm{d}t e^{i\omega t}\langle \hat{N}_D(t)\hat{N}_D(0)\rangle_H. \tag{3}$$

As usual, ω contains an infinitesimal imaginary part to ensure convergence of the integral. We will calculate $G(\omega)$ using $\hat{H}_T$ as perturbation. A substantial part of the calculations for the FM and the AF setup are identical, and we will distinguish between the two only when necessary.

The lowest-order contribution to $G(\omega)$ is quadratic in $\hat{H}_T$. It can be written as

$$G(\omega) = -\frac{i}{2}\int_0^\infty \mathrm{d}t e^{i\omega t}\int_{-\infty}^\infty \mathrm{d}t_1 \int_{-\infty}^\infty \mathrm{d}t_2 \times \langle T_t \hat{H}'_T(t_1)\hat{H}'_T(t_2)\hat{N}'_D(t)\hat{N}'_D(0)\rangle_0, \tag{4}$$

where $\hat{H}_T$ should be written in terms of $d_i^{(\dagger)}$'s and $r_i^{(\dagger)}$'s. The prime denotes an operator in the Heisenberg representation with respect to H_0. Since the operator $\hat{N}'_D(t)$ is defined such that $\hat{N}'_D(t)|\mathrm{gs}\rangle = 0$, where $|\mathrm{gs}\rangle$ denotes the ground state of the system under H_0, it follows immediately that there is only one time ordering that gives nonzero contributions. This time ordering leads to two different contributions to $G(\omega)$ that differ in the number of magnetic excitations in the intermediate state (either -1 or 1). After performing the integrations in Eq. (4) as well as a transformation to momentum space we obtain to second order in $\hat{H}_T$

$$G(\omega) = \frac{J_T^2}{4}\sum_{k,q}\left[\frac{\langle d_k d_k^\dagger\rangle_0\langle r_q^\dagger r_q\rangle_0}{(\varepsilon_k-\varepsilon_q)^2(\varepsilon_k-\varepsilon_q-\omega)} + \frac{\langle d_k^\dagger d_k\rangle_0\langle r_q r_q^\dagger\rangle_0}{(\varepsilon_q-\varepsilon_k)^2(\varepsilon_q-\varepsilon_k-\omega)}\right], \tag{5}$$

which is valid both for AF and FM systems. Equation (5) with $\omega = 0$ gives us C_M. When supplied with the relevant expectation values below, Eq. (5) tells us that the imaginary part of $G_{\mathrm{ret}}(\omega)$ at small $\omega \ll E_D^+, E_D^-$ [which determines R_M, see Eq. (1)] is zero to second order in $\hat{H}_T$. Hence, we need to analyze higher-order contributions to determine R_M. We will determine these contributions before analyzing Eq. (5) in more detail.

We have explicitly checked that the only time ordering in the fourth-order expression for $G(\omega)$ that leads to an imaginary contribution at small $\omega \ll E_D^+, E_D^-$ is given by $\langle \hat{H}'_T(t_1)\hat{N}'_D(t)\hat{H}'_T(t_2)\hat{H}'_T(t_3)\hat{N}'_D(0)\hat{H}'_T(t_4)\rangle_0$. This leads to six unique terms that cannot be excluded *a priori* and differ in the number of magnetic excitations in the intermediate states. We illustrate the procedure followed to determine $G(\omega)$ by focusing on the term for which the excess number of magnetic excitations on the dot varies as $0\to1\to0\to1\to0$. After performing the integrations over time as well as a transformation to momentum space we find the following contribution to $\mathrm{Im}[G(\omega)]$ at small ω to fourth order in $\hat{H}_T$ due to this term:

$$\left(\frac{J_T}{2}\right)^4 \sum_{k,\bar{k},q,\bar{q}} \frac{\langle d_k^\dagger d_k d_{\bar{k}}^\dagger d_{\bar{k}}\rangle_0\langle r_q r_{\bar{q}}^\dagger r_{\bar{q}} r_q^\dagger\rangle_0}{(\varepsilon_q-\varepsilon_k)^2(\varepsilon_q-\varepsilon_{\bar{k}})^2}\pi\delta(\varepsilon_q-\varepsilon_{\bar{q}}+\omega), \tag{6}$$

where it is understood that we need to take the continuum limit on the reservoir in order for the delta function to be well defined. The other terms that make up $G(\omega)$ can be calculated analogously, and we will refrain from repeating the required steps here.

Up to this point, our results for $G(\omega)$ are identical for FM and AF systems. The sole difference between the two arises now from the fact that $\langle d_k^\dagger d_{\bar{k}}\rangle_0$ and $\langle r_k^\dagger r_{\bar{k}}\rangle_0$ are different for AF and FM systems.

AF case.—Assuming $T\to0$, we put $\langle r_k^\dagger r_{\bar{k}}\rangle_0 = \langle d_k^\dagger d_{\bar{k}}\rangle_0 = \delta_{k,\bar{k}}\theta(\varepsilon_k-\mu_M)$. We can perform the summation over q in Eqs. (5)–(6) by replacing $\sum_q \to \nu_R\int \mathrm{d}q$, where $\nu_R = a|\partial\varepsilon_{q,R}/\partial q|^{-1}$ is the density of states in the reservoir. This leads to

$$C_M = t_{\mathrm{sd}}\sum_k \frac{1}{(\varepsilon_k-\mu_M)^2} \quad \text{and} \quad R_M = \frac{h}{2(g\mu_B)^2}, \tag{7}$$

where $t_{\mathrm{sd}} = (g\mu_B)^2\nu_R(J_T/2)^2$ is the "magnetic transparency" of a small magnetic dot. We note that C_M is well defined since $|\varepsilon_k-\mu_M| \geq \min[E_D^+, E_D^-]$. For large dots, we can also perform the summation over k using the density of states in the dot, ν_D (keeping in mind the previously discussed difficulties in the experimental realization of a finite $E_D^\pm$ for such dots). This leads to

$$C_M = t_{\mathrm{lD}}\left[\frac{1}{E_D^+}-\frac{1}{E_D^-}\right] \quad \text{and} \quad R_M = \frac{h}{(g\mu_B)^2}, \tag{8}$$

where $t_{\mathrm{lD}} = (g\mu_B)^2\nu_D\nu_R(J_T/2)^2$. In both cases we recover the fact that the spin resistance is universal in the sense that it does not depend on any microscopic parameters of the dot, not even on the coupling between the dot and the chain. This result is related to the fact that the AF spin-1/2 chain

maps onto one-dimensional fermions with interactions (see Supplemental Material [34]). This model has been studied extensively in the past few years in the context of electronic *RC* circuits [19–27]. The mapping allows us to support that the universality of the spin resistance remains valid for any value of the tunneling coupling [23,34].

Furthermore, the effect of electron-electron interactions in the reservoir and dot on the parameters of the *RC* circuit have been analyzed using a Luttinger liquid description [23,24]. Since these interactions translate to an $S_i^z S_{i+1}^z$ term in the spin chains, these results allow us to extend Eqs. (7)–(8) to finite values of $\Delta_{D(R)} \leq 1$. A finite value of $\Delta_{D(R)}$ corresponds to a deviation of the Luttinger liquid parameter from the noninteracting value $K = 1$. This simply leads to an additional factor $1/K$ in the result for R_M, so the resistance remains universal. The value of the capacitance C_M is changed in a nontrivial manner [23,24].

FM case.—The zero temperature limit is pathological for FM systems, since the Bose-Einstein distribution diverges at $T = 0$. Therefore, we will consider finite (but small) temperatures. Specifically, we will assume that $k_B T \ll E_D^+$, so that we can put $\langle d_k^\dagger d_{\bar{k}} \rangle_0 = 0$. Furthermore, we linearize $\langle r_q^\dagger r_{\bar{q}} \rangle_0 = n_B(E_R^+ + \epsilon_q)$ around the minimal value of the energy spectrum E_R^+. Substitution of these values into Eqs. (5)–(6) yields

$$C_M = \frac{t_{\mathrm{ID}}}{2} \frac{n_B(E_R^+)}{E_D^+ - E_R^+} \frac{\delta}{E_D^+ - E_R^+},$$
$$R_M = -\frac{2h}{(g\mu_B)^2} \left(\frac{E_D^+ - E_R^+}{\delta} \right)^2, \tag{9}$$

where $\delta = |n_B(E_R^+)/n_B'(E_R^+)|$. We find thus that the spin resistance is no longer universal for FM systems. The fact that the relaxation resistance is negative is not a fundamental issue but only means that the dynamical response of the spin capacitor is out of phase with the perturbation.

Note that the universality of the relaxation resistance in electronic interacting systems was shown in Ref. [23] to be intimately related to the Korringa-Shiba relation [35] (see also Garst *et al.* [36] for an extended version of this relation) which relates the imaginary part of the charge susceptibility to the square of its real part. Such a relation applies at or near a Fermi-liquid fixed point. It is therefore not surprising to find a nonuniversal behavior for R_M in the FM case, where low-energy excitations are bosonic.

We turn now to the possibility of using the setup displayed in Fig. 1 as a source of magnetic quasiparticles. Fève *et al.* have shown [16] experimentally that the electronic quantum *RC* circuit can be used as an on-demand single-electron source. Since the AF system maps onto the electronic system, we propose that the AF system can be used as an on-demand single-spinon source, with the simple substitution $eV_D(t) \to g\mu_B B_D(t)$. For FM systems, the situation is fundamentally different; since bosonic statistics allows for more than one magnon per momentum state, it is not possible to create a single-magnon source using the same mechanism. However, an on-demand few-magnon source appears feasible.

Finally, we comment on the possibility of measuring the properties of the magnetic quantum *RC* circuit experimentally. The ultimate implementation uses parallel spin chains, such as depicted in Fig. 1(a). Furthermore, molecular magnets [37–42] could be a good candidate to take the role of magnetic dot due to their beneficial properties, such as their increased size, chemical engineerability, and the possibility to control the spin state using electric instead of magnetic fields [43]. However, based on the magnitude of the spin currents [44] and the involved time scales, the implementation of the magnetic *RC* circuit in the above systems appears challenging, albeit not impossible. Therefore, we propose two alternative systems to test our predictions initially.

Our first proposal concerns ultracold atoms trapped in optical lattices. It has been shown [32,33] that these can be used to implement effective spin-1/2 Heisenberg chains (both AF and FM) with tunable anisotropy $\Delta_{R(D)}$. Furthermore, it is now possible [45–47] to measure the spin state in such systems dynamically, locally, and with single-spin precision.

We assume that the (effective) magnetic field $B_D(t)$ has the form $B_D(t) = B_D \cos(\omega_0 t)$. The validity of our results then requires $B_D, \hbar\omega_0 \ll |E_D^\pm|$. From Eq. (1) it follows that the resulting magnetization $M_D(t) = M_D^0 \cos(\omega_0 t) - M_D^1 \sin(\omega_0 t)$, where $M_D^0 = C_M B_D$ and $M_D^1 = R_M C_M^2 \omega_0 B_D$. To measure R_M, C_M, one has to be able to distinguish between these two contributions. For simplicity, we will give numbers for AF systems and small dots. Using Eq. (7) and taking the continuum limit on the dot, we estimate $M_D^0 \sim \xi(g\mu_B)$ and $M_D^1 \sim \xi^2(\hbar\omega_0/g\mu_B B_D)(g\mu_B)$, where $\xi = (J_T/J_D)(J_T/J_R)(g\mu_B B_D/E_D^+)$ is a small fraction. We have assumed that $|E_D^+| = |E_D^-|$ and $\varepsilon_{k,D(R)} = -J_{D(R)} \cos(ka)$. If we assume $J_T \lesssim J_D \approx J_R$ [see the discussion below Eq. (8) for the validity of this approximation] and $\hbar\omega_0 \approx g\mu_B B_D = 0.1 E_D^+$, it follows that a collection of $\sim 10^2$ parallel chains suffices to determine R_M, C_M in repeated measurements. For the smallest magnetic dots, $E_D^+ \approx J_D$. A representative value [33] for J_D is $J_D/\hbar \sim 0.1$ kHz, which leads to $\omega_0 = 10$ Hz, smaller than the typical lifetime of excitations in such systems [47].

Alternatively, artificially engineered atomic-spin chains could be used as sources of single spin excitations. We propose to use a setup very similar to that used in Ref. [6]. In our proposed setup, a single chain (consisting, for example, of several artificially coupled Fe atoms) acts as a magnetic dot, and is exchange coupled to a magnetic reservoir. By applying a pulsed magnetic field to the chain, a single magnetic excitation can be transported from the reservoir onto the chain or vice versa. Since it is possible to determine the spin state of the atoms in the chain using spin-polarized STM tips [6], the extra (or missing) spinon in the atomic spin chain can be detected. The magnitude of the applied magnetic field should be on the order of the charging energy $E_D^\pm \approx J_D a/L \approx 100$ mT for a spin chain consisting of 10

spins and $J_D \approx 1$ T. The typical tunneling time of the magnetic excitation can be estimated using Fermi's golden rule as $\tau^{-1} \approx J_T^2/(hE_D^{\pm})$. To illustrate, we find $\tau \approx 40$ ns for $J_T = 10$ mT and $\tau \approx 4$ μs for $J_T = 1$ mT.

Conclusion.—We have studied the magnetic *RC* circuit and computed its spin capacitance and resistance. We have shown that the resistance is universal for AF but not for FM systems. We have shown that our predictions can be presently tested with time-resolved experiments in ultracold atoms in optical lattices. This opens the path towards the realization of on-demand single spin-excitation emitters that would be one of the key ingredients for spintronics.

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