Cavity spintronics recently heralded non-local magnonic signal transfer between magnetic samples. Here we show that by including superconductors in the cavity, we can make use of these principles to bring superconducting spintronics to the macroscale. We analyze how a superconductor’s a.c. conductivity influences the spin dynamics of a spatially separated magnet, and we discuss the potential impact on spintronic applications.

The field of superconducting spintronics has been gathering pace in the last decade as the promise of achieving low dissipation spin and charge transport has been increasingly refined and realised [1–3]. The intriguing features rely on the proximity effect, whereby properties of one material can persist in an adjacent thin film. This places a tight nanometer constraint on the operational range in most cases. The most anomalously long-ranged persistence of superconductive signatures is reportedly up to the micrometer-range [4–5]. However, in this paper we highlight the untapped potential of superconducting spintronics to make use of advances in cavitrónics, and that photon-mediated processes can enable the detection of centimeter-ranged superconductive signatures. We provide a readily accessible example and discuss multiple directions for exploration to highlight the potential for innovation in superconducting spintronic applications.

Cavity spintronics, or cavitronics, is an emerging interdisciplinary field in which microwave or optical cavity photon modes can couple to magnons (also called spin waves). Experiments have shown strong coupling of cavity modes to both ferri- and ferromagnets [6–7]. This is observed as a hybridization of the photon and magnon modes, indicated by avoided crossings/Rabi splitting in the normal mode frequency spectrum. It was recently shown that magnonic interactions between two non-local magnetic samples can be mediated by the cavity modes [8–10]. That is, the cavity photons permit coupling between spatially separated magnetic samples, and consequently facilitates the transmission of spintronic signatures on macroscopic length scales. We explore the question of magnons coupling non-locally to excitations in a superconductor.

Light with frequencies above the superconducting gap breaks Cooper pairs and thus weakens the superconductivity. However, light can also enhance or induce superconductivity [11–13]. In-cavity manipulation of a superconductive component might appear restrictive, demanding effective screening of the contact wires while maintaining the quality factor of the cavity, but also this has been achieved experimentally recently [14]. In that case, researchers succeeded in driving a black box transmon qubit inside a cavity, coupling the oscillations between the two levels of the qubit to the microwave cavity modes. The transmon qubit is engineered by using the non-linearity of a superconducting Josephson junction to create an effective two-level system, as in circuit quantum electrodynamics [15]. In this way, the qubit-cavity coupling generated quite some excitement about the potential prospect of unifying quantum optics and solid state qubit quantum computing [16–17].

Qubit-cavity coupling demonstrated that it is feasible to successfully screen wiring to a superconducting system inside microwave cavities. However, the superconductivity in that case is used as a means to generate a two-level system, i.e., realize a qubit, and not as a means to probe and use the superconductive signatures themselves. In this article, we will argue that there is considerable potential to do just that.

We begin by considering the setup illustrated in Fig. 1.

**FIG. 1.** (Color online) The proposed model for inducing macroscale photon-mediated superconducting signatures in a magnet (not to scale). The photonic microwave resonator of dimensions \(d_x, d_y, d_z\) contains a thin and small superconducting wire segment (SC) along the \(y\)-direction and with a cross-sectional area \(A\), connected to an alternating current (a.c.) source via screened wiring through the cavity walls, as well as a small ferromagnetic sphere (FM) with a uniform magnetization \(m\). The FM and the SC are positioned at \(r_{FM}\) and \(r_{SC}\) respectively, corresponding to extrema of the magnetic and electric components of the cavity mode \(B_{cav}\) and \(E_{cav}\). Across the SC, \(E_{cav}\) is directed along the \(y\)-axis, and across the FM, \(B_{cav}\) is directed along the \(x\)-axis. The FM is additionally subjected to a strong external magnetostatic field \(B_{ext}\) such that \(|B_{ext}| > |B_{cav}|\), which fixes the precessional axis of \(m\) along the \(z\)-direction. In our set-up, the selected cavity mode is the TE_{201} mode. The SC current, cavity mode and FM mode couple resonantly at the input a.c. frequency \(\omega\). The relative amounts of supercurrent and resistive currents passed through the SC is modulated by the temperature \(T\).
picting a microwave cavity at temperature $T$. The cavity contains an electrically screened thin wire, which has a small exposed superconducting segment (SC), connected to an alternating current (a.c.) source, as well as a small ferromagnetic sphere (FM). The internal current density $J$ and electric field $E_{SC}$ of the SC are treated as uniform; i.e., internal spatial variations are neglected. The SC and the FM are placed in regions of maximum electric and magnetic fields $E_{cav}$ and $B_{cav}$ of a selected cavity mode, respectively. The dimensions of the SC and the FM are assumed sufficiently small for the local fields across their respective regions to be approximately uniform, and their spatial extension are effectively taken to be line-like and point-like at positions $r_{SC}$ and $r_{FM}$, respectively.

The SC is directed along the $y$-direction, and has a critical temperature $T_c$. The a.c. source produces a signal of frequency $\omega$, which is resonant with the cavity frequency and the frequency of the precessing FM magnetization. By lowering the temperature of the cavity, we pass through the superconducting transition and induce a change in the superconductors conductivity. This in turn alters the excitation of the cavity, and the resultant effect on the spin dynamics in the magnet can then be harnessed as a non-local detector. That is, by exploiting the mutually resonant coupling to the cavity, it is possible to probe the superconducting transition via a change in the magnonic precession response. For the sake of simplicity in the following illustrative example, we neglect the cavity back-action on the superconductor, and consider only weak coupling here.

To provide a concrete example, we consider the TE$_{201}$ cavity mode. In this case, $E_{cav}$ is directed along the $y$-axis over the SC, and $B_{cav}$ is directed along the $z$-axis over the FM. $B_{cav}$ then couples predominantly to the Kittel mode of the FM, i.e. the uniform mode of the spherical spin field, quantified by the unit magnetization vector $m$. The FM is additionally exposed to a relatively strong external magnetic field $B_{ext}$ such that $|B_{ext}| \gg |B_{cav}|$, which fixes the precessional axis of $m$ along the $z$-direction. $|B_{ext}|$ additionally regulates the resonance frequency of the spin field mode, and reduces to small perturbations the impact of $B_{cav}$ on the motion of $m$.

The resonance frequency of the TE$_{201}$ mode is determined by a given set of dimensional parameters $\{d_x, d_z\}$ of the cavity, and one may thus match the resonance frequency of the Kittel mode and the frequency of the input a.c. to this, by appropriately adjusting $|B_{ext}|$ and $\omega$.

The current response of a superconductor to an applied electric field, taking into account both frequency and temperature dependency, may be derived from microscopic theories of superconductivity, such as BCS theory or Eliashberg theory. Mattis–Bardeen theory is derived from the former $^{[19]}$, and provides accurate descriptions of the optical conductivity of BCS superconductors. However, these theories are generally cumbersome to deal with analytically, and will in this paper be reserved for numerical calculations. To analytically model the transition from resistive to superconducting current in the SC, we may instead employ the well established framework of the Drude model-based two-fluid model $^{[21]}$.

The SC is treated as two parallel channels carrying normal ($n$) and superconducting ($s$) electrons, respectively. The superconducting channel is characterized by an asymptotically infinite relaxation time $\tau_s \rightarrow \infty$, and for the normal channel, a low input frequency $\omega \tau_n \ll 1$ relative to the relaxation time of the normal electrons is assumed. By the Drude model, the relationship between the current density responses and $E_{SC}$ are thus

$$\frac{dJ_s(\omega, T, t)}{dt} = \frac{N_s(T)e^2}{m_e}E_{SC}(\omega, T, t), \quad (1)$$

$$\frac{J_n(\omega, T, t)}{\tau_n} = \frac{N_n(T)e^2}{m_e}E_{SC}(\omega, T, t), \quad (2)$$

where $m_e$ is the electron mass, and $J_s$ and $N_s$ are the current and electron densities of the respective channels. It is then clear that for sinusoidal time dependencies there is a relative phase difference of $\pm \pi/2$ between the contributions of $J_s$ and $J_n$ to $E_{SC}$ in a current-driven system. $E_{SC}$ thus acquires a phase relative to the net current density $J = J_s + J_n$ between 0 and $\pm \pi/2$. We argue that this phase shift can be used to bridge superconducting and spintronic circuits via non-local coupling to magnons. In this case it can monitor the superconducting transition, and be implemented as a superconducting switch. More broadly, it opens the door for wider investigations of macro-scale effects in superconducting circuits.

Upon connecting the SC to an a.c. current source, the magnitude of the net current density is $J(\omega, t) = I \exp(i\omega t)/A$, where $I$ is the current amplitude, $A$ is the cross-sectional area of the SC, and $\omega$ is the input frequency. Inserting this into Eqs. (1) and (2), one finds that

$$E_{SC}(\omega, T, t) = \frac{I}{A\sigma(\omega, T)} \exp(i\omega t),$$

where $\sigma(\omega, T)$ is the real (as $\omega \rightarrow 0$) and negative imaginary part of the SC conductivity $\sigma = \sigma_1 - i\sigma_2$, as a function of temperature $T$, for various frequency inputs $\omega$. Material parameters for Nb are used ($T_c = 9.26$ K) $^{[13]}$. $\sigma_0$ is the normal state direct current conductivity. These plots are generated numerically using Mattis–Bardeen theory $^{[19]}$.

![FIG. 2. (Color online) Intersections of the real ($\sigma_1$) and negative imaginary ($\sigma_2$) part of the SC conductivity $\sigma = \sigma_1 - i\sigma_2$ as a function of temperature $T$, for various frequency inputs $\omega$. Material parameters for Nb are used ($T_c = 9.26$ K) $^{[13]}$. $\sigma_0$ is the normal state direct current conductivity. These plots are generated numerically using Mattis–Bardeen theory $^{[19]}$.](image-url)
where
\[ \sigma(\omega, T) = \frac{e^2}{m_e} \left( N_n(T) \tau_n - i \frac{N_s(T)}{\omega} \right) = \sigma_1(T) - i \sigma_2(\omega, T). \] (4)

The temperature dependency of \( N_s \) and by extension \( \sigma_1 \) and \( \sigma_2 \), is phenomenologically taken to be
\[ N_s(T) = N \left[ 1 - \left( \frac{T}{T_c} \right)^4 \right], \quad N_n(T) = N \left( \frac{T}{T_c} \right)^4, \] (5)
where \( N \) is the total density of electrons, \( T_c \) is the critical temperature of the SC, and \( T \leq T_c \). For the purpose of analytic insight we retain this simple form, although we include the standard temperature modification of the gap in the numerics based on the Mattis–Bardeen theory. Above \( T_c \), \( \sigma \) reduces to the normal metal direct current conductivity \( \sigma_0 = N c^2 \tau_n/m_e \). Note that according to Mattis–Bardeen theory, \( \sigma_1 \) is frequency dependent; near \( T_c \), it has a pronounced coherence peak at lower frequencies, and a kink at higher frequencies due to optical excitations across the superconducting gap, as illustrated in Fig. 2 [19, 20]. Neither of these features are captured by the two-fluid model. Nevertheless, in terms of the relative magnitudes of \( \sigma_1 \) and \( \sigma_2 \), and their point of intersection marking the boundary between the superconducting and the resistive regime, the two-fluid model and Mattis–Bardeen theory coincide very well at the experimentally relevant lower frequencies. Fig. 2 thus shows the predicted temperatures for the transition between normal and superconducting current [22].

\( E_{SC} \) and \( E_{cav} \) are assumed to be purely tangential to the SC–cavity interface in our set-up (see Fig. 1). Thus, by the continuity of the tangential electric field across any interface, \( E_{SC}(r_{SC}, \omega, T, t) = E_{cav}(r_{SC}, \omega, T, t) \) at the surface of the SC. Upon computing the cavity modes by imposing rectangular boundary conditions on the fields, e.g. as done in Ref. [23], one then finds that across the FM and specifically for the TE\(_{201}\) mode, \( B_{cav} \) at the FM is [24]
\[ B_{cav}(r_{FM}, \omega, T, t) = \hat{B}_{cav}(r_{FM}, \omega, T, t) = -\frac{\pi E_{cav}(r_{SC}, \omega, T, t)}{i \omega d_z} \hat{x}. \] (6)

Furthermore, the resonance frequency of the TE\(_{201}\) mode is
\[ \omega = c \sqrt{\left( \frac{2 \pi}{d_x} \right)^2 + \left( \frac{\pi}{d_z} \right)^2}. \] (7)

With \( d_x \) and \( d_z \) given, the resonance frequency is equal to the input a.c. frequency by appropriate tuning of \( \omega \).

The precessional motion of the FM magnetization vector \( \mathbf{m} \) is adequately described by the Landau–Lifshitz–Gilbert (LLG) equation:
\[ \frac{\partial \mathbf{m}(\omega, T, t)}{\partial t} = -\gamma \mathbf{m}(\omega, T, t) \times \mathbf{B}(\omega, T, t) + \alpha \mathbf{m}(\omega, T, t) \times \frac{\partial \mathbf{m}(\omega, T, t)}{\partial t}. \] (8)

Here, \( \gamma \) and \( \alpha \) are the gyromagnetic ratio and the phenomenological damping parameter of the LLG equation, respectively, and \( \mathbf{B} \) is the effective magnetic field inside the FM, including the external field, the demagnetization field and the magnetocrystalline anisotropy field [25, 26]. The latter two are generally influenced by the geometry and crystal structure of the FM, and may influence the resonance frequency and orbit of \( \mathbf{m} \). We assume that the FM has an easy axis such as \((111)\) for YIG [27], and that this axis coincides with the \( z \)-direction. Assuming furthermore that the magnetostatic field \( B_{ext} \) across the FM is much greater than the demagnetization and anisotropy fields, the influence of the latter two may be neglected. This is reasonably expected to hold down to an input frequency of 5 GHz [27, 30]. The effective magnetic field across the FM is then
\[ B(\omega, T, t) = B_{cav}(r_{FM}, \omega, T, t) + B_{ext} \hat{z} \]
\[ = \frac{\pi E_{cav}(r_{SC}, \omega, T, t)}{i \omega d_z} \hat{x} + B_{ext} \hat{z}. \] (9)

Under the assumption that \( |B_{cav}| \gg |B_{ext}| \), the \( z \)-component of \( \mathbf{m} \) may be taken to be unity to first order in the magnitude of the consequently small precessing component, and in Eq. (9), terms of higher order than linear in \( B_{cav} \), or the remaining components \( m_x \) and \( m_y \) of \( \mathbf{m} \), may be neglected. In addition, the coupling between the cavity mode and the FM is taken to be resonant. Under these conditions, the precessional motion of \( \mathbf{m} \) is elliptical; furthermore, one may then safely proceed to solve the LLG equation with complex time dependencies \( \exp(i \omega t) \) in \( \mathbf{B} \) and \( \mathbf{m} \), and finally extract the real parts of these quantities as the physical solutions. The expression for \( \mathbf{m} \) therefore has the general linearized form
\[ \mathbf{m}(\omega, T, t) \approx \hat{z} + m_p(\omega, T, t), \] where the precessing compo-

**FIG. 3.** (Color online) Phase of the magnon precession \( \varphi_m \), as given by Eq. (11), as a function of temperature \( T \) and for frequency inputs \( \omega \), using Mattis–Bardeen theory to compute the SC conductivity. Material parameters for Nb and YIG are used, with \( T_c = 9.26 \, K \) and \( \alpha = 10^{-5} [18, 26, 31, 32] \). This phase may be measured relative to the input signal passed through the SC, and its value indicates the relative presence of supercurrent and resistive current in the SC. For this case, \( \varphi_m \approx \varphi_m - \pi/2 \).
ternal magnetic field to the input a.c. frequency by an appropriate tuning of the ex-

Reinserting the solutions for the magnitude of the precessing component \(|m_p|\) as given by Eq. (10), as a function of temperature \(T\) and for frequency inputs \(\omega\), using Mattis–Bardeen theory to compute the SC conductivity. Material parameters for Nb, a microwave cavity and YIG are used, with \(T_c = 9.26\, \text{K}\), \(\alpha = 10^{-3}\), \(\gamma = 176\, \text{GHz/T}\). \(I = 0.6\, \text{A}, A = 10^{-11}\, \text{cm}^2\) and \(d_z = 5\, \text{cm}\) [18, 26, 31, 32]. Within experimental limits such as the critical current of the SC, the decrease in magnitude for increasing critical current of the SC, the decrease in magnitude for increasing frequencies may be compensated for by increasing the input current. Furthermore, to ensure consistency with deposited platinum layer. As discussed in the text, the analytics would then require the inclusion of the demagnetization field and associated shift in resonance, but it would not otherwise alter the physics. Further, one may include the self-consistent back-action on the superconductor and perform a readout on the SC side.

The work shown that photon-mediated superconducting signatures are a feasible way to provide a bridging circuit for spintronic applications. In device design this can feature as a superconductive switch, but also to monitor the superconducting transition and critical temperature of the superconductor directly.

However, the importance of the result also goes beyond these applications as it opens up a plethora of interesting investigatory avenues. For example, by switching from a conventional singlet superconductor to a triplet source (either intrinsically \(p\)-wave or odd-frequency \(s\)-wave), then there are no longer two simple coupling relationships to the cavity as in the case of the a.c.-driven oscillators in Eqs. (1) and (2). The nature of this coupling remains to be explored, but it seems plausible in that case that one may employ the cavity setup to probe and differentiate between the different current compo-

\[
\varphi_{m_x}(\omega, T, t) \equiv \arg(m_x(\omega, T)) \\
\approx \arctan\left(\frac{\sigma_2(\omega, T)}{\sigma_1(T)} + \frac{\alpha}{2}\right),
\]

\[
\varphi_{m_y}(\omega, T) \equiv \arg(m_y(\omega, T)) \\
\approx \arctan\left(\frac{\sigma_2(\omega, T)}{\sigma_1(T)} - \frac{\alpha}{2} - \frac{\pi}{2}\right).
\]

Reinserting the solutions for \(m_x\) and \(m_y\) into Eq. (8), then taking the absolute value of both sides, yields \(\omega = |\gamma| B_{\text{ext}}\), equal to the input a.c. frequency by an appropriate tuning of the external magnetic field \(B_{\text{ext}}\). Furthermore, to ensure consistency with linear response and elliptical orbits, we should have that the magnitude of the precessing component \(|\text{Re}(m_p)|\) \(\ll 1\) [33]. \(\text{Re}(m_p)\) relates to the cone angle of the precession, and is given by

\[
|\text{Re}(m_p(\omega, T, t))| \approx \\
|m_y(\omega, T)| \sqrt{2\alpha \cos^2\left(\omega t + \varphi_{m_y} + \frac{\alpha}{2} - \frac{\pi}{2} - \theta\right) - \alpha + 1},
\]

where

\[
|m_y(\omega, T)| \approx \frac{|\gamma| \pi I}{2A|\sigma(\omega, T)| \omega^2 d_z \alpha},
\]

\[
\theta \approx \frac{3\pi + \alpha}{4}.
\]
ments. This may make the super cavitronics setup an interesting new tool for probing unconventional superconductors.

For the physical picture presented above, it is sufficient to consider a classical description of the coupling. However, it would be interesting to explore a microscopic picture along the line of cavity quantum electrodynamics as outlined in Ref. [40]. In that case we can of course not neglect the details of the mesoscopic circuit by tracing over the mesoscopic degrees of freedom, meaning the mathematical approach becomes rather involved. Nevertheless, it is expected to yield valuable insight to the case of fermionic reservoirs in a cavity.

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