Optical Cooling of Magnons

Sharma, Sanchar; Blanter, Yaroslav M.; Bauer, Gerrit E.W.

DOI
10.1103/PhysRevLett.121.087205

Publication date
2018

Published in
Physical Review Letters

Citation (APA)
Sharma, S., Blanter, Y. M., & Bauer, G. E. W. (2018). Optical Cooling of Magnons. *Physical Review Letters*, *121*(8), Article 087205. https://doi.org/10.1103/PhysRevLett.121.087205

Optical Cooling of Magnons

Sanchar Sharma,[1] Yaroslav M. Blanter,[1] and Gerrit E. W. Bauer[2,1]
[1]*Kavli Institute of NanoScience, Delft University of Technology, 2628 CJ Delft, The Netherlands*
[2]*Institute for Materials Research & WPI-AIMR & CSRN, Tohoku University, Sendai 980-8577, Japan*

(Received 10 April 2018; published 22 August 2018)

Inelastic scattering of light by spin waves generates an energy flow between the light and magnetization fields, a process that can be enhanced and controlled by concentrating the light in magneto-optical resonators. Here, we model the cooling of a sphere made of a magnetic insulator, such as yttrium iron garnet, using a monochromatic laser source. When the magnon lifetimes are much larger than the optical ones, we can treat the latter as a Markovian bath for magnons. The steady-state magnons are canonically distributed with a temperature that is controlled by the light intensity. We predict that such a cooling process can significantly reduce the temperature of the magnetic order within current technology.

DOI: 10.1103/PhysRevLett.121.087205

A great achievement of modern physics is the Doppler cooling of trapped atoms by optical lasers [1,2] down to temperatures of micro-Kelvin [3]. Subsequently, even macroscopic mechanical objects, such as membranes and cantilevers, have been cooled to their quantum mechanical ground state [4–8] by blueshifting the stimulated emission using an optical cavity [4,5]. Cavity optomechanics is a vibrant field that achieved successful Heisenberg uncertainty-limited mechanical measurements, the generation of entangled light-mechanical states, and ultrasensitive gravitational wave detection [8]. An optical cryocooler based on solid state samples [9] can be superior due to its compactness and lack of moving components [10]. Optical cooling has been demonstrated for glass [9,11] and envisioned for semiconductors [10,12,13].

An analogous cooling of a magnet would generate interesting opportunities. Magnetization couples to microwaves [14–19], electric currents [18,20,21], mechanical motion [17,22–24], and, indeed, light [25,26]. Spin waves are the elementary excitations of the ferromagnetic order, which are quantized as bosonic magnons. Similar to phonons, magnons may be considered noninteracting up to relatively high temperatures and are Planck distributed at thermal equilibrium. However, there are important differences as well: Magnons have mass and chirality [27,28], both of which are tunable by an external static magnetic field. Their long wavelength dispersion in thin films is highly anisotropic, with minima in certain directions that can collect the Bose-Einstein condensate of magnons [29–31]. Magnons can be used as quantum transducers between microwaves and optical light [19,32] or between superconducting and flying qubits [33].

Motivated by the potential of a ferromagnet as a versatile quantum interface at low temperatures, we discuss here the potential of optical cooling of magnons. The magnon-photon interaction gives rise to inelastic Brillouin light scattering (BLS) [34], which is a well-established tool to study magnon dispersion and dynamics [25,35,36]. Recently, several groups carried out BLS experiments on spheres made of ferrimagnetic insulator yttrium iron garnet (YIG) [37–42], which has a very high magnetic quality factor ($\sim 10^5$) [43–45] and supports ferromagnetic magnons with long coherence times ($\sim\mu$s) [33,46,47]. YIG spheres are commercially available for microwave applications, but are also good infrared light cavities due to their large refractive index and high optical quality [48–50], making them good optomagnonic resonators [37–41, 51–55]. Via proximity optical fibers or prisms, external laser light can efficiently excite whispering gallery modes (WGMs), i.e., the optical modes circulating in extremal orbits of dielectric spheroids [56,57].

Early BLS experiments on YIG spheres discovered an asymmetry in the redshifted (Stokes) and blueshifted (anti-Stokes) sidebands [25,58] due to an interference of Faraday and Cotton-Mouton effects. When more photons are scattered into the blue sideband, light effectively extracts energy from the magnons and in principle cools them. This asymmetry is enhanced by 3–4 orders of magnitude in a WGM cavity [38–41] due to resonant enhancement of the scattering cross section [40,53,54]. The asymmetry can be controlled by the polarization and wave vector of the light. Optomagnonic scattering is enhanced for a triple resonance condition [40,59–62] by tuning both the input and the scattered photon frequency to the optical resonances of the cavity. In contrast, optomechanical cooling [4,5,8] requires detuning the input laser from a cavity resonance with correspondingly reduced scattering and cooling rate. In this Letter, we predict that modern technology and materials can significantly reduce the temperature of the magnetic order, showing the potential to manipulate magnons using light.

We derive below rate equations for photons and magnons to estimate the steady-state magnon number that can be

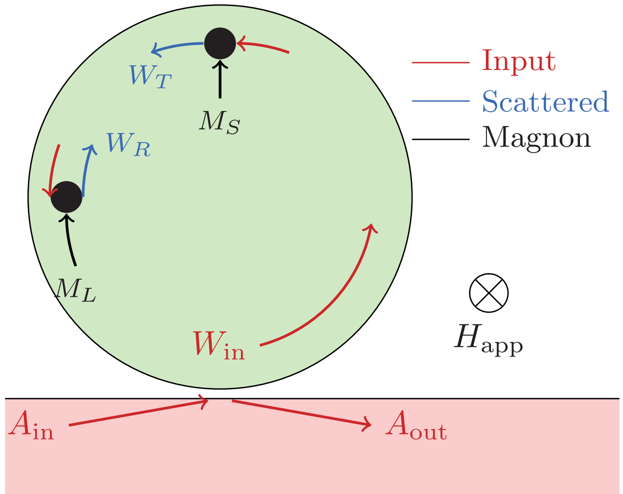

FIG. 1. Optomagnonic cooling setup: A ferromagnetic sphere in contact with an optical waveguide. A magnetic field $H_{\rm app}$ (normal to the waveguide) is applied to saturate the magnetization. Input light with amplitude $A_{\rm in}$ is evanescently coupled to a WGM $W_{\rm in}$. We focus on anti-Stokes scattering by two types of magnons that are characterized by their angular momentum [53]. A small angular momentum magnon M_S maintains the direction of WGMs, converting $W_{\rm in}$ to W_T. $W_{\rm in}$ can be reflected into W_R by absorbing a large angular momentum magnon M_L. Theoretically, both the cases can be treated in the same formalism.

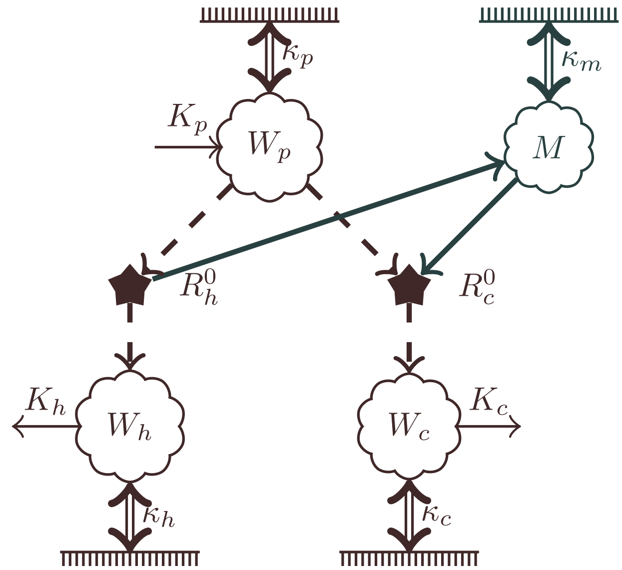

FIG. 2. Light-induced cooling of a magnon, M. A proximity fiber or prism is coupled to the WGMs W_x with a coupling constant K_x, exciting W_p while collecting the scattered W_c and W_h. The photons are inelastically scattered by the magnon $W_p + M \to W_c$ and $W_p \to W_h + M$ at single particle rates R_c^0 and R_h^0 respectively, derived in the text. All modes are coupled to their respective thermal baths by leakage rates κ_x. When κ_c is much larger than the corresponding scattering rate, the bath associated with W_c can become an efficient channel for dissipation of the magnons in M.

reached as a function of material and device parameters. We consider a spherical magnetic insulator with a high index of refraction that is transparent at the input light frequency (Fig. 1) and magnetization perpendicular to the WGM orbits that are excited by proximity coupling to an external laser. We single out two groups of magnon modes that couple preferentially to the WGMs [53]: The small angular momentum (including the Kittel) magnons, M_S in Fig. 1, and the large angular momentum magnons, the chiral Damon-Eshbach modes M_L. The theory presented below is valid for both types of magnons.

We can understand the basic physics by the minimal model sketched in Fig. 2. We focus on a single incident WGM W_p with index p and frequency ω_p. It is occupied by [8]

$$n_p = \frac{4K_p}{(\kappa_p + K_p)^2} \frac{P_{\rm in}}{\hbar\omega_p} \tag{1}$$

photons, with κ_p being the intrinsic linewidth, K_p the leakage rate into the proximity coupler, and $P_{\rm in}$ the input light power. An optically active magnon M (with either small or large angular momentum) is annihilated $W_p + M \to W_c$ or created $W_p \to W_h + M$ by BLS, where W_c and W_h are blue-shifted and redshifted sideband WGMs, respectively.

We first derive a simple semiclassical rate equation for the nonequilibrium steady-state magnon number, $n_m^{(sc)}$ (the superscript distinguishes the estimate from n_m as more rigorously derived below). The thermal bath absorbs and injects magnons at rates $\kappa_m n_m^{(sc)}(n_{\rm th}+1)$ and $\kappa_m n_{\rm th}(n_m^{(sc)}+1)$ respectively, where κ_m is the inverse magnon lifetime and $n_{\rm th}$ is the occupation of the bath at magnon frequency ω_m. At an ambient temperature T,

$$n_{\rm th} = \left[\exp\left(\frac{\hbar\omega_m}{k_B T}\right) - 1\right]^{-1}. \tag{2}$$

The optical cooling rate is $R_c^0 n_p n_m^{(sc)}$, where R_c^0 is the anti-Stokes scattering rate of one W_p photon by one M magnon and we assumed that there are no photons in W_c. The latter is justified because of small optomagnonic couplings compared to WGM dissipation rates, $\sim 2\pi \times 0.1$–1 GHz [38–40] while $R_c^0 n_p n_m^{(sc)}$ is at most $\sim\kappa_m \sim 2\pi \times 1$ MHz. In the absence of dissipation, Fermi's golden rule gives $R_c^0 = 2\pi |g_c|^2 \delta(\omega_p + \omega_m - \omega_c)$, where $\hbar g_c$ is the optomagnonic coupling and $\{\omega_p, \omega_c, \omega_m\}$ are the frequencies of $\{W_p, W_c, M\}$, respectively. When W_c has a finite lifetime, the δ function is broadened into a Lorentzian, giving

$$R_c^0 = \frac{|g_c|^2 (\kappa_c + K_c)}{(\omega_p + \omega_m - \omega_c)^2 + (\kappa_c + K_c)^2/4}, \tag{3}$$

where κ_c is its intrinsic linewidth, and K_c is its leakage rate into the proximity coupler. Similarly, the optical heating rate is $R_h^0 n_p (n_m^{(sc)} + 1)$, where R_h^0 is given by Eq. (3) with $g_c, \omega_c, \kappa_c \to g_h, \omega_h, \kappa_h$ and $\omega_m \to -\omega_m$. In deriving $R_{c,h}^0$, we ignore the magnon linewidth since $\kappa_m \ll \kappa_c, \kappa_h$ [40,53]. In the steady state the cooling and heating rates are equal, leading to the estimate

$$n_m^{(sc)} = \frac{\kappa_m n_{\text{th}} + R_h^0 n_p}{\kappa_m + (R_c^0 - R_h^0) n_p}. \tag{4}$$

This agrees with the result from the more precise theory discussed below, thus capturing the essential processes correctly (*a posteriori*). However, the rate equation cannot access noise properties beyond the magnon number that are important for thermodynamic applications. Further, it does not differentiate between a coherent precession of the magnetization and the thermal magnon cloud with the same number of magnons.

In order to model the cooling process more rigorously, we proceed from a model Hamiltonian for a system with three photon and one magnon modes. In the Hamiltonian $\hat{H}_S = \hat{H}_0 + \hat{H}_{\text{OM}}$ [53]

$$\hat{H}_0 = \hbar\omega_p \hat{a}_p^\dagger \hat{a}_p + \hbar\omega_c \hat{a}_c^\dagger \hat{a}_c + \hbar\omega_h \hat{a}_h^\dagger \hat{a}_h + \hbar\omega_m \hat{m}^\dagger \hat{m}, \tag{5}$$

and $\hat{a}_x$ and $\hat{m}$ are the annihilation operators for photon W_x with $x \in \{p, c, h\}$ and magnon M. The optomagnonic coupling in the rotating wave approximation reads [53]

$$\hat{H}_{\text{OM}} = \hbar g_c \hat{a}_p \hat{a}_c^\dagger \hat{m} + \hbar g_h \hat{a}_p \hat{a}_h^\dagger \hat{m}^\dagger + \text{H.c.}, \tag{6}$$

where g_c and g_h are the scattering amplitudes and H.c. is the Hermitian conjugate.

In the rotating frame of the "envelope" operators $\hat{W}_x(t) \triangleq \hat{a}_x(t) e^{i\omega_x t}$ and $\hat{M}(t) \triangleq \hat{m}(t) e^{i\omega_m t}$ the (Heisenberg) equation of motion for $\hat{M}$ becomes [4,63]

$$\dot{\hat{M}} = -i g_h \hat{W}_p \hat{W}_h^\dagger e^{i\delta_h t} - i g_c^* \hat{W}_p^\dagger \hat{W}_c e^{-i\delta_c t} - \frac{\kappa_m}{2} \hat{M} - \sqrt{\kappa_m} \hat{b}_m, \tag{7}$$

where $\delta_h = \omega_h + \omega_m - \omega_p$ and $\delta_c = \omega_c - \omega_m - \omega_p$ are the detunings from the scattering resonances. $\hat{b}_m(t)$ is the stochastic magnetic field generated by the interaction of M with phonons [64] and/or other magnons [65], whose precise form depends on the microscopic details of the interaction [66].

We assume that the correlators of $\hat{b}_m$ obey the fluctuation-dissipation theorem for thermal equilibrium [67,68]. When $\kappa_m \ll k_B T/\hbar$, which is satisfied for $\kappa_m \sim 2\pi \times 1$ MHz [38–40] and $T \gg 50$ μK, the (narrow band filtered) noise is effectively white and generates a canonical Gibbs distribution of the magnons in steady state [63]. Their statistics are $\langle \hat{b}_m(t) \rangle = 0$, $\langle \hat{b}_m^\dagger(t') \hat{b}_m(t) \rangle = n_{\text{th}} \delta(t - t')$ and $\langle \hat{b}_m(t') \hat{b}_m^\dagger(t) \rangle = (n_{\text{th}} + 1) \delta(t - t')$, where n_{th} is defined in Eq. (2).

For weak scattering relative to the input power, we can ignore any back-action on W_p such that its dynamics are governed only by the proximity coupling. When W_p is in a coherent state, $\langle \hat{W}_p(t) \rangle = \sqrt{n_p}$ and $\langle \hat{W}_p^\dagger(t') \hat{W}_p(t) \rangle = n_p$, where n_p is given by Eq. (1).

The photons in W_c are generated by the optomagnonic coupling and dissipated into their thermal bath, with the Heisenberg equation of motion [4,63]

$$\frac{d\hat{W}_c}{dt} = -i g_c \hat{W}_p \hat{M} e^{i\delta_c t} - \frac{\kappa_c + K_c}{2} \hat{W}_c - \sqrt{\kappa_c} \hat{b}_c - \sqrt{K_c} \hat{A}_c, \tag{8}$$

where $\hat{b}_c$ and $\hat{A}_c$ are noise operators. The physical origins of $\hat{b}_c$ and finite lifetime κ_c^{-1} are scattering by impurities, surface roughness, and lattice vibrations. K_c is the leakage rate of W_c into the proximity coupler and $\hat{A}_c$ is the vacuum noise from the latter into W_c. The noise sources are white for sufficiently small κ_c. $\langle \hat{X}_c(t) \rangle = 0$, $\langle \hat{X}_c^\dagger(t') \hat{X}_c(t) \rangle = 0$ and $\langle \hat{X}_c(t') \hat{X}_c^\dagger(t) \rangle = \delta(t - t')$ where $X \in \{\hat{b}_c, \hat{A}_c\}$, because the thermal occupation of photons at infrared and visible frequencies is negligibly small at room temperature $e^{-\hbar\omega_c/(k_B T)} \approx 0$.

The solution to Eq. (8) is $\hat{W}_c(t) = \hat{W}_{c,\text{th}}(t) + \hat{W}_{c,\text{OM}}(t)$. The thermal contribution is

$$\hat{W}_{c,\text{th}} = \int_0^t e^{-(\kappa_c + K_c)(t-\tau)/2} [-\sqrt{\kappa_c} \hat{b}_c(\tau) - \sqrt{K_c} \hat{A}_c(\tau)] d\tau \tag{9}$$

where the origin of time is arbitrary. For $t, t' \to \infty$, we get the equilibrium statistics $\langle \hat{W}_{c,\text{th}}^\dagger(t') \hat{W}_{c,\text{th}}(t) \rangle = 0$ and

$$\langle \hat{W}_{c,\text{th}}(t') \hat{W}_{c,\text{th}}^\dagger(t) \rangle = \exp\left(-\frac{(\kappa_c + K_c)|t - t'|}{2}\right), \tag{10}$$

independent of the optomagnonic coupling. $\hat{W}_{c,\text{OM}}$ can be simplified by the adiabaticity of the magnetization dynamics that follows from $\kappa_m \ll \kappa_c$. When $\hat{M}$ is treated as a slowly varying constant,

$$\hat{W}_{c,\text{OM}}(t) \approx -i g_c \hat{M}(t) \int_0^t e^{-(\kappa_c + K_c)(t-\tau)/2} \hat{W}_p(\tau) e^{i\delta_c \tau} d\tau. \tag{11}$$

$\hat{W}_h(t)$ is obtained by the substitution $c \to h$ and $\hat{M} \to \hat{M}^\dagger$ in Eqs. (9)–(11).

We can now rewrite Eq. (7) as

$$\frac{d\hat{M}}{dt} = -\left(\frac{\kappa_m}{2}\hat{M} + \sqrt{\kappa_m}\hat{b}_m\right) + \hat{O}_c + \hat{O}_h. \quad (12)$$

with cooling and heating operators that reflect the light-scattering processes in Fig. 2:

$$\hat{O}_c = \hat{\mathcal{N}}_c + i\hat{\Sigma}_c\hat{M}, \quad (13)$$

$$\hat{O}_h = -\hat{\mathcal{N}}_h^\dagger + i\hat{\Sigma}_h^\dagger\hat{M}. \quad (14)$$

Focusing on the cooling process, we distinguish the self-energy,

$$\hat{\Sigma}_c = i|g_c|^2\int_0^t e^{[i\delta_c+(\kappa_c+K_c)/2](\tau-t)}\hat{W}_p^\dagger(t)\hat{W}_p(\tau)d\tau, \quad (15)$$

from the noise operator,

$$\hat{\mathcal{N}}_c(t) = -ig_c^*\hat{W}_p^\dagger(t)\hat{W}_{c,\mathrm{th}}(t)e^{-i\delta_c t}. \quad (16)$$

In the weak-coupling regime we may adopt a mean-field approximation by replacing $\hat{\Sigma}_c$ by its average,

$$\langle\hat{\Sigma}_c\rangle = -\bar{\omega}_c + i\frac{\bar{\kappa}_c}{2} \triangleq \frac{|g_c|^2 n_p}{\delta_c - i(\kappa_c + K_c)/2}, \quad (17)$$

where $\bar{\omega}_c$ is the (reactive) shift of the magnon resonance and $\bar{\kappa}_c$ the optical contribution to the magnon linewidth.

The noise $\hat{\mathcal{N}}_c$ can be interpreted as the vacuum fluctuations of W_c entering the magnon subsystem via the optomagnonic interaction. $\hat{\mathcal{N}}_c$ has a very short correlation time $\sim(\kappa_c + K_c)^{-1}$ [see Eq. (10)] compared to magnon dynamics $\sim\kappa_m^{-1}$, and thus can be treated as a white noise source with $\langle\hat{\mathcal{N}}_c(t)\rangle = 0$, $\langle\hat{\mathcal{N}}_c^\dagger(t)\hat{\mathcal{N}}_c(t')\rangle = 0$, and $\langle\hat{\mathcal{N}}_c(t')\hat{\mathcal{N}}_c^\dagger(t)\rangle \approx V_c\delta(t-t')$. By integrating over time and using the correlation functions of $\hat{W}_p$ and $\hat{W}_{c,\mathrm{th}}$

$$V_c = \frac{4|g_c|^2 n_p(\kappa_c + K_c)}{4\delta_c^2 + (\kappa_c + K_c)^2} = \bar{\kappa}_c, \quad (18)$$

defined in Eq. (17). $\bar{\kappa}_c/\kappa_m$ at resonance $\delta_c = 0$ is the cooperativity between the magnons and W_c photons due to the coupling mediated by W_p photons.

Analogous results hold for $\hat{O}_h$, with substitutions $c \to h$ in Eqs. (15)–(18). We arrive at

$$\frac{d\hat{M}}{dt} \approx -i(\bar{\omega}_c + \bar{\omega}_h)\hat{M} - \frac{\kappa_{\mathrm{tot}}}{2}\hat{M} - \sqrt{\kappa_{\mathrm{tot}}}\hat{b}_{\mathrm{tot}}, \quad (19)$$

where $\kappa_{\mathrm{tot}} = \kappa_m + \bar{\kappa}_c - \bar{\kappa}_h$ and $\sqrt{\kappa_{\mathrm{tot}}}\hat{b}_{\mathrm{tot}} = \sqrt{\kappa_m}\hat{b}_m - \hat{\mathcal{N}}_c + \hat{\mathcal{N}}_h^\dagger$. The fluctuations of the total noise follow from Eq. (18)

$$\langle\hat{b}_{\mathrm{tot}}^\dagger(t')\hat{b}_{\mathrm{tot}}(t)\rangle \approx n_m\delta(t-t'), \quad (20)$$

$$\langle\hat{b}_{\mathrm{tot}}(t')\hat{b}_{\mathrm{tot}}^\dagger(t)\rangle \approx (n_m+1)\delta(t-t'), \quad (21)$$

where

$$n_m = \frac{\kappa_m n_{\mathrm{th}} + \bar{\kappa}_h}{\kappa_m + \bar{\kappa}_c - \bar{\kappa}_h}. \quad (22)$$

Equation (19) is equivalent to the equation of motion for magnons in equilibrium after the substitutions $\omega_m \to \omega_m + \bar{\omega}_c + \bar{\omega}_h$, $\kappa_m \to \kappa_{\mathrm{tot}}$, and $n_{\mathrm{th}} \to n_m$. It implies that the magnons in the nonequilibrium steady state are still canonically distributed with the density matrix

$$\hat{\rho}_{\mathrm{ne}} = \exp\left(\frac{-\hbar\omega_m\hat{n}_m}{k_B T_{\mathrm{ne}}}\right)\left(\mathrm{Tr}\left[\exp\left(\frac{-\hbar\omega_m\hat{n}_m}{k_B T_{\mathrm{ne}}}\right)\right]\right)^{-1} \quad (23)$$

where the number operator $\hat{n}_m = \hat{m}^\dagger\hat{m}$ and the nonequilibrium magnon temperature T_{ne} is implicitly defined by Eq. (22) and

$$n_m = \left[\exp\left(\frac{\hbar\omega_m}{k_B T_{\mathrm{ne}}}\right) - 1\right]^{-1}. \quad (24)$$

We get $\langle\hat{M}_x\rangle = \langle\hat{M}_y\rangle = 0$, which implies that light scattering does not induce a coherent magnon precession, in contrast to a resonant ac magnetic field. n_m is the average number of magnons that can be larger or smaller than the equilibrium value n_{th}. The result is consistent with $n_m^{(sc)}$ [see Eq. (4)] because $\bar{\kappa}_{c,h} = R_{c,h}^0 n_p$ as expected from Fermi's golden rule. The canonical distribution implies that the steady-state magnon entropy is maximized for the given number of magnons, n_m.

When $\bar{\kappa}_h - \bar{\kappa}_c > \kappa_m$, that is, when heating by the laser overcomes the intrinsic magnon damping, the system becomes unstable, leading to runaway magnon generation and self-oscillations [52,69,70]. The instability is regularized by magnon-magnon scattering, not included in our theory.

Here we focus on the cooling scenario in which $\bar{\kappa}_h \ll \bar{\kappa}_c$ [53]. Magnon cooling can be monitored by the intensity of the blueshifted sideband. Using the input-output formalism [63,71], the scattered light amplitude in the rotating frame is

$$\hat{A}_{\mathrm{out}}(t) = -\sqrt{K_c}\hat{W}_c(t). \quad (25)$$

It can be converted into the output power by $P_{\mathrm{out}} = \hbar\omega_c\langle\hat{A}_{\mathrm{out}}^\dagger(t)\hat{A}_{\mathrm{out}}(t)\rangle$, which is independent of time in steady state. With impedance matching, $\kappa_{p,c} = K_{p,c}$, and at the triple resonance, $\delta_c = 0$,

$$\frac{P_{\mathrm{out}}}{P_{\mathrm{in}}} = \frac{|g_c|^2}{\kappa_c\kappa_p}\frac{\kappa_m n_{\mathrm{th}}}{\kappa_m + 2|g_c|^2 n_p/\kappa_c} \propto \frac{1}{1 + P_{\mathrm{in}}/P_s}, \quad (26)$$

defining the saturation power

$$P_s \triangleq \frac{\hbar\omega_p\kappa_p\kappa_c\kappa_m}{2|g_c|^2}. \quad (27)$$

To leading order $P_{\rm out} \propto P_{\rm in}$ [39,53], but saturates when the magnon number becomes small, which is experimental evidence for magnon cooling. P_s is the input power that halves the number of magnons.

Equation (23) is the reduced density matrix of a single magnon mode, M. At long wavelengths and small magnon numbers, the magnon-magnon interactions may be disregarded, so each magnon mode can be treated independently. The total $\hat{\rho}$ is a direct product of the reduced density matrices of the form (23), where ω_m and $T_{\rm ne}$ depend on the mode index.

For a YIG sphere with parameters $\omega_c \approx \omega_p = 2\pi \times 300$ THz (free space wavelength 1 μm), an optical Q factor $\omega_p/(2\kappa_p) = \omega_c/(2\kappa_c) = 10^6$, [39], magnon linewidth $\kappa_m = 2\pi \times 1$ MHz, and optomagnonic coupling $g_c = 2\pi \times 10$ Hz [53], we get $P_s = 140$W. Trying to match this with $P_{\rm in}$ is not useful since laser-induced lattice heating [10] will overwhelm the cooling effect. However, P_s can be significantly reduced by large magnon-photon coupling. Doping YIG with bismuth can increase g_c tenfold [50], bringing P_s down to $\sim$1 W. The spatial overlap between the magnons and photons [53] can be engineered in ellipsoidal or nanostructured magnets [72] which can increase g_c further by an order of magnitude, giving $P_s \sim 10$ mW. For an ambient temperature $T = 1$ K and magnon frequency $\omega_m = 2\pi \times 10$ GHz, the thermal magnon number $n_{\rm th} = 1.62$. For $P_{\rm in} = \{P_s/20, P_s, 5P_s\}$ the steady-state magnon numbers are $n_m = \{1.55, 0.81, 0.27\}$ and temperatures $T_{ne} = \{0.96, 0.60, 0.31\}$ K respectively. At an optimistic $P_s = 10$ mW, the above input power corresponds to $n_p = \{3 \times 10^6, 5 \times 10^7, 3 \times 10^8\}$ intracavity photons respectively. Cooling is experimentally observable for relatively small powers $P_{\rm in} < P_s/20$, which should be achievable by optimizing the magnon-photon coupling.

In summary, we estimate the cooling power due to BLS of light by magnons in an optomagnonic cavity. Because of the large mismatch of optical and magnonic timescales, the photon degree of freedom can be eliminated by renormalizing the magnon frequency and damping, cf. Eq. (19), and a light-controlled effective temperature Eq. (22). Current technology and materials are close to achieving significant cooling of magnons, envisioning the possibility of light-controlled magnon manipulation.

We thank J. Haigh, M. Elyasi, K. Satoh, K. Usami, and A. Gloppe for helpful input and discussions. This work is financially supported by the Nederlandse Organisatie voor Wetenschappelijk Onderzoek (NWO) as well as JSPS KAKENHI Grant No. 26103006.

[1] T. Hänsch and A. Schawlow, Opt. Commun. **13**, 68 (1975).
[2] W. D. Phillips, Rev. Mod. Phys. **70**, 721 (1998).
[3] P. D. Lett, R. N. Watts, C. I. Westbrook, W. D. Phillips, P. L. Gould, and H. J. Metcalf, Phys. Rev. Lett. **61**, 169 (1988).
[4] F. Marquardt, J. P. Chen, A. A. Clerk, and S. M. Girvin, Phys. Rev. Lett. **99**, 093902 (2007).
[5] I. Wilson-Rae, N. Nooshi, W. Zwerger, and T. J. Kippenberg, Phys. Rev. Lett. **99**, 093901 (2007).
[6] A. D. O'Connell, M. Hofheinz, M. Ansmann, R. C. Bialczak, M. Lenander, E. Lucero, M. Neeley, D. Sank, H. Wang, M. Weides, J. Wenner, J. M. Martinis, and A. N. Cleland,, Nature (London) **464**, 697 (2010).
[7] J. Chan, T. P. M. Alegre, A. H. Safavi-Naeini, J. T. Hill, A. Krause, S. Groblacher, M. Aspelmeyer, and O. Painter, Nature (London) **478**, 89 (2011).
[8] M. Aspelmeyer, T. J. Kippenberg, and F. Marquardt, Rev. Mod. Phys. **86**, 1391 (2014).
[9] B. C. Edwards, J. E. Anderson, R. I. Epstein, G. L. Mills, and A. J. Mord, J. Appl. Phys. **86**, 6489 (1999).
[10] M. Sheik-Bahae and R. Epstein, Laser Photonics Rev. **3**, 67 (2009).
[11] R. I. Epstein, M. I. Buchwald, B. C. Edwards, T. R. Gosnell, and C. E. Mungan, Nature (London) **377**, 500 (1995).
[12] M. Sheik-Bahae and R. I. Epstein, Phys. Rev. Lett. **92**, 247403 (2004).
[13] K. Usami, A. Naesby, T. Bagci, B. Melholt Nielsen, J. Liu, S. Stobbe, P. Lodahl, and E. S. Polzik, Nat. Phys. **8**, 168 (2012).
[14] Ö. O. Soykal and M. E. Flatté, Phys. Rev. Lett. **104**, 077202 (2010).
[15] X. Zhang, C.-L. Zou, L. Jiang, and H. X. Tang, Phys. Rev. Lett. **113**, 156401 (2014).
[16] Y. Tabuchi, S. Ishino, T. Ishikawa, R. Yamazaki, K. Usami, and Y. Nakamura, Phys. Rev. Lett. **113**, 083603 (2014).
[17] X. Zhang, C.-L. Zou, L. Jiang, and H. X. Tang, Sci. Adv. **2**, e1501286 (2016).
[18] L. Bai, M. Harder, Y. P. Chen, X. Fan, J. Q. Xiao, and C.-M. Hu, Phys. Rev. Lett. **114**, 227201 (2015).
[19] M. Kostylev and A. Stashkevich, arXiv:1712.04304.
[20] Y. Kajiwara, K. Harii, S. Takahashi, J. Ohe, K. Uchida, M. Mizuguchi, H. Umezawa, H. Kawai, K. Ando, K. Takanashi, S. Maekawa, and E. Saitoh,, Nature (London) **464**, 262 (2010).
[21] A. V. Chumak, V. I. Vasyuchka, A. A. Serga, and B. Hillebrands, Nat. Phys. **11**, 453 (2015).
[22] C. Kittel, Phys. Rev. **110**, 836 (1958).
[23] A. Kamra, H. Keshtgar, P. Yan, and G. E. W. Bauer, Phys. Rev. B **91**, 104409 (2015).
[24] T. Kikkawa, K. Shen, B. Flebus, R. A. Duine, K.-i. Uchida, Z. Qiu, G. E. W. Bauer, and E. Saitoh, Phys. Rev. Lett. **117**, 207203 (2016).
[25] A. Borovik-Romanov and N. Kreines, Phys. Rep. **81**, 351 (1982).
[26] F. Hansteen, A. Kimel, A. Kirilyuk, and T. Rasing, Phys. Rev. B **73**, 014421 (2006).
[27] R. Damon and J. Eshbach, J. Phys. Chem. Solids **19**, 308 (1961).
[28] D. D. Stancil and A. Prabhakar, *Spin Waves: Theory and Applications* (Springer, New York, 2008), pp. 169–201.
[29] S. O. Demokritov, V. E. Demidov, O. Dzyapko, G. A. Melkov, A. A. Serga, B. Hillebrands, and A. N. Slavin, Nature (London) **443**, 430 (2006).
[30] S. M. Rezende, Phys. Rev. B **79**, 174411 (2009).

[31] N. P. Proukakis, D. W. Snoke, and P. B. Littlewood, *Universal Themes of Bose-Einstein Condensation* (Cambridge University Press, New York, 2017).
[32] R. Hisatomi, A. Osada, Y. Tabuchi, T. Ishikawa, A. Noguchi, R. Yamazaki, K. Usami, and Y. Nakamura, Phys. Rev. B **93**, 174427 (2016).
[33] Y. Tabuchi, S. Ishino, A. Noguchi, T. Ishikawa, R. Yamazaki, K. Usami, and Y. Nakamura, Science **349**, 405 (2015).
[34] D. D. Stancil and A. Prabhakar, *Spin Waves: Theory and Applications* (Springer, New York, 2008) pp. 223–260.
[35] T. Sebastian, K. Schultheiss, B. Obry, B. Hillebrands, H. Schultheiss, and B. Obry, Front. Phys. **3** (2015).
[36] S. Klingler, H. Maier-Flaig, R. Gross, C.-M. Hu, H. Huebl, S. T. B. Goennenwein, and M. Weiler, Appl. Phys. Lett. **109**, 072402 (2016).
[37] J. A. Haigh, S. Langenfeld, N. J. Lambert, J. J. Baumberg, A. J. Ramsay, A. Nunnenkamp, and A. J. Ferguson, Phys. Rev. A **92**, 063845 (2015).
[38] A. Osada, R. Hisatomi, A. Noguchi, Y. Tabuchi, R. Yamazaki, K. Usami, M. Sadgrove, R. Yalla, M. Nomura, and Y. Nakamura, Phys. Rev. Lett. **116**, 223601 (2016).
[39] X. Zhang, N. Zhu, C.-L. Zou, and H. X. Tang, Phys. Rev. Lett. **117**, 123605 (2016).
[40] J. A. Haigh, A. Nunnenkamp, A. J. Ramsay, and A. J. Ferguson, Phys. Rev. Lett. **117**, 133602 (2016).
[41] A. Osada, A. Gloppe, R. Hisatomi, A. Noguchi, R. Yamazaki, M. Nomura, Y. Nakamura, and K. Usami, Phys. Rev. Lett. **120**, 133602 (2018).
[42] J. A. Haigh, N. J. Lambert, S. Sharma, Y. M. Blanter, G. E. W. Bauer, and A. J. Ramsay, Phys. Rev. B **97**, 214423 (2018).
[43] V. Cherepanov, I. Kolokolov, and V. L'vov, Phys. Rep. **229**, 81 (1993).
[44] A. A. Serga, A. V. Chumak, and B. Hillebrands, J. Phys. D **43**, 264002 (2010).
[45] M. Wu and A. Hoffmann, in *Recent Advances in Magnetic Insulators—From Spintronics to Microwave Applications*, edited by M. Wu and A. Hoffmann, Solid State Physics Vol. 64 (Academic Press, Cambridge, MA, 2013), pp. 1–392.
[46] X. Zhang, C.-L. Zou, N. Zhu, F. Marquardt, L. Jiang, and H. X. Tang, Nat. Commun. **6**, 8914 (2015).
[47] Y. Tabuchi, S. Ishino, A. Noguchi, T. Ishikawa, R. Yamazaki, K. Usami, and Y. Nakamura, C.R. Phys. **17**, 729 (2016).
[48] D. L. Wood and J. P. Remeika, J. Appl. Phys. **38**, 1038 (1967).
[49] W. Wettling, J. Magn. Magn. Mater. **3**, 147 (1976).
[50] D. Lacklison, G. Scott, H. Ralph, and J. Page, IEEE Trans. Magn. **9**, 457 (1973).
[51] T. Liu, X. Zhang, H. X. Tang, and M. E. Flatté, Phys. Rev. B **94**, 060405 (2016).
[52] S. Viola-Kusminskiy, H. X. Tang, and F. Marquardt, Phys. Rev. A **94**, 033821 (2016).
[53] S. Sharma, Y. M. Blanter, and G. E. W. Bauer, Phys. Rev. B **96**, 094412 (2017).
[54] A. Osada, A. Gloppe, Y. Nakamura, and K. Usami, arXiv:1711.09321.
[55] J. Graf, H. Pfeifer, F. Marquardt, and S. V. Kusminskiy, arXiv:1806.06727.
[56] A. N. Oraevsky, Quantum Electron. **32**, 377 (2002).
[57] M. R. Foreman, J. D. Swaim, and F. Vollmer, Adv. Opt. Photonics **7**, 168 (2015).
[58] W. Wettling, M. G. Cottam, and J. R. Sandercock, J. Phys. C **8**, 211 (1975).
[59] J. M. Dobrindt and T. J. Kippenberg, Phys. Rev. Lett. **104**, 033901 (2010).
[60] A. H. Safavi-Naeini and O. Painter, New J. Phys. **13**, 013017 (2011).
[61] A. Mari and J. Eisert, Phys. Rev. Lett. **108**, 120602 (2012).
[62] G. Bahl, M. Tomes, F. Marquardt, and T. Carmon, Nat. Phys. **8**, 203 (2012).
[63] C. W. Gardiner and M. J. Collett, Phys. Rev. A **31**, 3761 (1985).
[64] C. Vittoria, S. D. Yoon, and A. Widom, Phys. Rev. B **81**, 014412 (2010).
[65] E. Schlömann, J. Phys. Chem. Solids **6**, 242 (1958).
[66] Á. Rivas and S. F. Huelga, *Open Quantum Systems: An Introduction*, Springer Briefs in Physics (Springer-Verlag, Berlin Heidelberg, 2012), pp. 1–97.
[67] W. F. Brown, Phys. Rev. **130**, 1677 (1963).
[68] L. Landau and E. Lifshitz, in *Statistical Physics*, 3rd ed., edited by L. Landau and E. Lifshitz (Butterworth-Heinemann, Oxford, 1980), pp. 333–400.
[69] H. Suhl, J. Phys. Chem. Solids **1**, 209 (1957).
[70] S. M. Rezende and F. M. de Aguiar, Proc. IEEE **78**, 893 (1990).
[71] A. A. Clerk, M. H. Devoret, S. M. Girvin, F. Marquardt, and R. J. Schoelkopf, Rev. Mod. Phys. **82**, 1155 (2010).
[72] F. Heyroth, C. Hauser, P. Trempler, P. Geyer, F. Syrowatka, R. Dreyer, S. G. Ebbinghaus, G. Woltersdorf, and G. Schmidt, arXiv:1802.03176.