

Hallmarks of Ballistic Terahertz Magnon Currents in an Antiferromagnetic Insulator

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Efficient transport of spin angular momentum is expected to play a crucial role in future spintronic devices, which will potentially operate at frequencies reaching the terahertz range. On the other hand, antiferromagnetic insulators exhibit significant potential for facilitating ultrafast pure spin currents by terahertz magnons. Consequently, we here use femtosecond laser pulses to trigger ultrafast spin currents across thin films of the model antiferromagnet NiO in Py|NiO|Pt stacks, where permalloy (Py) and Pt serve as spin-current source and detector, respectively. We find that the spin-current pulses traversing NiO reach a velocity up to 40 nm/ps and experience increasing attenuation, delay, and broadening as the NiO thickness is increased. We can consistently explain our observations by ballistic transport of incoherent magnons. The relaxation length of 20 nm is 1 order of magnitude larger than for ballistic conduction-electron spin transport in metals and, thus, relevant for potential applications. Methodologically, our all-optical and contact-free spin-conductance approach has large potential to characterize terahertz magnon transport in insulators with any kind of magnetic order.

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Introduction—The primary objective of spintronics research is to augment electronic devices, which are based on the charge of the electron, by the spin of the electron. In magnetic insulators, pure spin currents carried by magnons, rather than spin-polarized conduction electrons, are particularly promising due to the strongly reduced Joule dissipation [1–3]. Antiferromagnetic insulators may even facilitate magnon transport in the terahertz (THz) range [4–6]. Characteristics, such as low magnetic damping, enhanced transverse spin correlation near the Néel temperature [7], and the superfluidity of Bose-Einstein condensates at low temperature [8,9], position antiferromagnetic insulators as natural candidates for efficient ultrafast spin transport.

However, the mechanism underlying magnon propagation in antiferromagnetic insulators remains a subject of intense debate, particularly concerning whether or not spins are transported coherently [7]. At frequencies below the magnon gap, spin currents injected via temperature gradients [10–13] or spin pumping [14–17] typically relax over distances of less than 10 nm. Nevertheless, spin transport over distances up to several micrometers was observed in antiferromagnetic insulators with weak net magnetization [18–21]. This remarkably long distance was attributed to coherent magnon propagation [22]. However, further experimental and theoretical investigation into spin transport in antiferromagnetic insulators is crucial to push transport to THz bandwidths.

To trigger THz spin transport, one can take advantage of the spin voltage that is induced in a ferromagnetic-metal (FM) layer through ultrafast optical excitation of the FM [Fig. 1(a)] [23]. This spin voltage was concluded to inject THz spin currents into metals, such as Cu [24,25], and insulators, such as MgO [26]. After propagation across the interlayer, the spin current can be detected in a subsequent heavy-metal (HM) layer by conversion into a transverse charge current that is monitored through the THz electromagnetic pulse it emits [23,27] [Fig. 1(a)]. This versatile technique can potentially be

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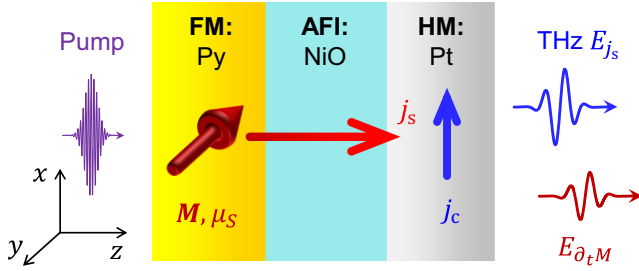


FIG. 1. Schematic for studying ultrafast spin transport through an antiferromagnetic-insulator layer AFI of NiO in a FM|AFI|HM stack. A femtosecond pump pulse induces a spin voltage $\mu_s(t)$ in the ferromagnetic-metal layer FM of permalloy (Py), which drives a spin current through the AFI. The spin current $j_s(t)$ arriving in the heavy-metal layer HM of Pt is converted into a charge current $j_c(t) \propto j_s(t)$ and monitored via the THz electric field $E_{j_s}(t)$ it emits. Ultrafast quenching of the FM magnetization $\mathbf{M}$ leads to the emission of an additional field $E_{\partial_t M}(t)$. An external magnetic field sets the direction of $\mathbf{M}$ along the y axis.

used to measure the spin conductance of an antiferromagnetic-insulator layer (AFI) by inserting it as an interlayer. Recent pioneering works have successfully reported spintronic THz emission from optically excited FM|AFI|HM stacks with an AFI of NiO [28,29]. However, important questions remain, for example, whether the spin current is driven by a gradient of spin voltage or temperature, akin to a Seebeck effect. Even more importantly, we need to reveal the dispersion, mean velocity, and relaxation length of the potential THz magnon current and evaluate whether it is predominantly carried by coherent or incoherent magnons.

In this Letter, we study the ultrafast dynamics of spin currents across thin films of the model antiferromagnetic insulator NiO by extending ultrabroadband THz spin-conductance spectroscopy [26] to a Py|NiO|Pt stack. Clear spin-current transients are observed for a NiO thickness up to 30 nm with a time resolution down to 120 fs. The ultrafast spin-transport dynamics are well described by a model based on ballistic transport of incoherent rather than coherent magnons. Our findings provide a direct visualization of spin transport mediated by THz short-wavelength magnons in antiferromagnets. Therefore, our results are relevant for both fundamental magnon studies and the development of ultrafast low-dissipation spintronic devices based on antiferromagnetic insulators.

Experimental concept—Figure 1 illustrates our THz spin-conductance method [26] for measuring ultrafast spin transport across a layer AFI within a FM|AFI|HM stack. The FM has magnetization $\mathbf{M} = M\mathbf{u}_y$ along the y axis with unit vector $\mathbf{u}_y$. A femtosecond laser pulse transiently heats the FM and generates a time-dependent spin voltage $\mu_s(t)$, i.e., an excess of spin angular momentum [23]. The latter drives a spin current with density $j_s(t)$ through the AFI toward the HM. Upon reaching the HM, the $j_s(t)$ is converted into an in-plane charge current $j_c(t)$ via the

inverse spin Hall effect, resulting in a THz electric field $E_{j_s}(t)$ directly behind the sample. In the far field, E_{j_s} is detected as an electro-optic signal $S_{j_s}(t)$.

By comparing $S_{j_s}(t)$ for a sample with an AFI of thickness d to that of a reference sample FMIHM without AFI ($d = 0$ nm), we extract the time-dependent spin conductance $G_{\text{AFI}}(t)$ of the AFI. Here, $G_{\text{AFI}}(t)$ denotes the spin current in response to an impulsive spin voltage $\mu_s(t) \propto \delta(t)$ [26].

Note that, in addition to magnon transport through the AFI, other contributions to magnetic-field-dependent THz emission need to be considered [30]. In particular, heating of the FM by the pump pulse can induce spin transfer to the crystal lattice. It, thus, changes the FM magnetization $\mathbf{M}$ at a rate $\partial_t \mathbf{M}(t) \propto \mu_s(t)$, leading to an additional electric field $E_{\partial_t M}(t)$ (Fig. 1) [23,31,32].

In our FM|AFI|HM|sub stacks, Py(3 nm) serves as the FM, NiO(d) as the AFI, and Pt(4 nm) as the HM on a MgO(500 μm) substrate (sub). More details on the samples and experimental setup are provided in Appendix A.

Raw data—Figure 2(a) shows typical THz electro-optic signals $S(t)$ odd in M [Eq. A1] from Py|NiO(d)|Pt for $d = 0, 3, 8, 18$, and 30 nm, as measured with a ZnTe(1 mm) detector. The strong THz signal from the reference sample Py|Pt (i.e., $d = 0$ nm) is rescaled by a factor of 1/20. The signals $S(t)$ for all $d > 0$ nm exhibit similar dynamics in the region $t < 1$ ps, in particular a similar amplitude at $t = 0$ ps. Importantly, the later dynamics at $t > 1$ ps depends sensitively on d up to 30 nm [Fig. 2(a)]. We confirm that the odd-in- M signal $S(t)$ is independent of both the pump polarization and sample azimuth. The associated THz electric field is perpendicular to $\mathbf{M} = M\mathbf{u}_y$, consistent with the contributions E_{j_s} and $E_{\partial_t M}$ indicated in Fig. 1.

Frequency domain—For the moment, let us assume that the total transient THz electric field E is dominated by the sum $E_{j_s} + E_{\partial_t M}$ (Fig. 1), where both E_{j_s} and $E_{\partial_t M}$ are driven by the transient spin voltage μ_s in the FM. Note that features of spin transport (j_s) and local magnetization quenching ($\partial_t M$) are strongly blurred in the measured electro-optic signal traces $S = S_{j_s} + S_{\partial_t M}$ in Fig. 2(a). This behavior arises from the noninstantaneous dynamics $\mu_s(t)$ of the spin voltage in the FM, the conversion of the charge current due to j_s and $\partial_t M$ into the electric field $E = E_{j_s} + E_{\partial_t M}$, the propagation of E to the detector and its electro-optic detection (Fig. 1) [23,26].

To isolate the spin-transport component from the total signal S , we compare it to the reference signal $S_{\text{ref}} = S|_{d=0}$ obtained from the control sample Py|Pt. Because S_{ref} is dominated by direct spin transport from the FM into the HM, it closely reflects the dynamics $\mu_s(t) \propto j_s(t)|_{d=0}$ of the spin voltage. As shown in Supplemental Material, Sec. 3 [33], the ratio of the Fourier-transformed signals,

$$\tilde{R}(\omega) = \frac{\tilde{S}(\omega)}{\tilde{S}_{\text{ref}}(\omega)} = \frac{\tilde{G}_{\text{AFI}}(\omega)}{g_{\text{ref}}} + b, \quad (1)$$

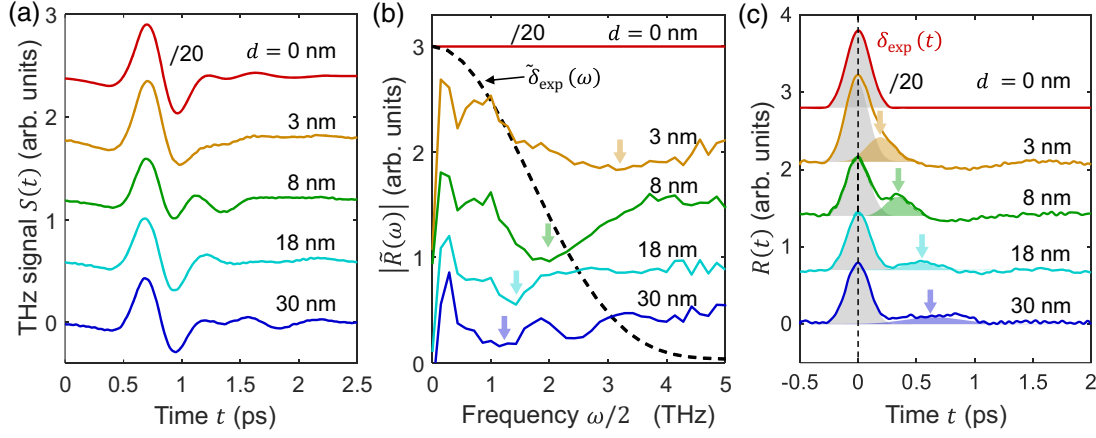


FIG. 2. THz-emission signals from optically excited Py|NiO(d)|Pt stacks. (a) Electro-optic signal $S(t)$ odd in M [Eq. (A1)] for various NiO thicknesses d , measured in ZnTe (1 mm). The signal from Py|Pt (i.e., $d=0$ nm) is the reference signal $S_{\text{ref}}(t)$. (b) Ratio $\tilde{R}(\omega)$ of the Fourier-amplitude spectra of $S(t)$ and $S_{\text{ref}}(t)$ [Eq. (1)]. Arrows indicate spectral dips. The filter function $\tilde{\delta}_{\text{exp}}(\omega)$ (dashed line) represents the finite bandwidth of our setup. (c) Signal $(R * \delta_{\text{exp}})(t)$, obtained by inverse Fourier transformation of $\tilde{R}(\omega)\tilde{\delta}_{\text{exp}}(\omega)$ from panel (b). The shaded areas at $t=0$ ps (dashed line) and at $t>0$ ps (arrows) indicate the two components making up $R(t)$. The function $\delta_{\text{exp}}(t)$ is the inversely Fourier-transformed $\tilde{\delta}_{\text{exp}}$ and represents the impulsive spin voltage inducing $(R * \delta_{\text{exp}})(t)$ [Eq. (1)]. In all panels, all curves are offset for better visibility, except $\tilde{\delta}_{\text{exp}}(\omega)$ and the signals for $d=30$ nm. All curves for $d=0$ nm are rescaled by a factor of $1/20$.

reveals the spin conductance $\tilde{G}_{\text{AFI}}(\omega) = \tilde{j}_s(\omega)/\tilde{\mu}_s(\omega)$ of the AFI in the frequency domain. The coefficients $1/g_{\text{ref}}$ and b quantify the respective weight of the contribution of j_s and $\partial_t M$ (Fig. 1). They are frequency independent because g_{ref} is the spin conductance of the Py/Pt interface, which is instantaneous [26], and because $(\partial_t M)(t)$ follows the driving force $\mu_s(t)$ instantaneously [23].

Figure 2(b) shows $|\tilde{R}(\omega)|$ for various NiO thicknesses d . As expected from Eq. (1), $|\tilde{R}(\omega)| = 1$ holds for $d=0$ nm. For $d>0$ nm, we observe the emergence of a spectral dip, whose width and center frequency both decrease systematically as d increases from 3 to 30 nm. This evolution of $\tilde{R}(\omega)$ with d indicates a substantial and thickness-dependent contribution of spin transport and, thus, the spin conductance $\tilde{G}_{\text{AFI}}(\omega)$ to the THz signal.

Time domain—To better understand the spectra in Fig. 2(b), we Fourier transform $\tilde{R}(\omega)$ into the time domain. According to Eq. (1), the resulting time-domain response is expected to follow

$$R(t) = \frac{G_{\text{AFI}}(t)}{g_{\text{ref}}} + b\delta(t). \quad (2)$$

Here, the time-domain spin conductance $G_{\text{AFI}}(t)$ corresponds to the spin current $j_s(t)$ that reaches the HM in response to an impulsive spin voltage $\mu_s(t) \propto \delta(t)$ in the FM [26]. Likewise, the term $b\delta(t)$ captures the contribution from the local magnetization dynamics $\partial_t M(t)$ driven by the same impulsive spin voltage.

Because the bandwidth of our experiment is finite, we approximate the ideal $\delta(t)$ by the peaklike function $\delta_{\text{exp}}(t)$ [Fig. 2(c)]. Consequently, we extract the convolution $R * \delta_{\text{exp}}$ [Eq. (B1)] rather than R (see Appendix B).

For all NiO thicknesses d , the $R * \delta_{\text{exp}}$ traces exhibit a prominent peak around $t=0$ ps [Fig. 2(c)]. Its shape agrees well with the driving force δ_{exp} , and its amplitude is approximately independent of d . We further find a second, delayed peak, whose position shifts monotonically toward later times as d increases. The behavior of R is fully compatible with Eq. (2). In particular, the δ -like peak at $t=0$ ps can be attributed to ultrafast local demagnetization ($\partial_t M \propto \mu_s = \delta_{\text{exp}}$), while the time-delayed peak corresponds to spin transport across the NiO layer, as described by $j_s = G_{\text{AFI}} * \mu_s = G_{\text{AFI}} * \delta_{\text{exp}}$. Further aspects of the extracted $R * \delta_{\text{exp}}$ and its connection to superluminal magnon propagation suggested previously [28] are discussed in Appendix C.

THz spin conductance—As our primary interest lies in ultrafast spin transport through NiO, we extract the spin conductance $G_{\text{AFI}}(t)$ by subtracting the initial $\delta_{\text{exp}}(t)$ peak from the $(R * \delta_{\text{exp}})(t)$ traces in Fig. 2(c) [see Eq. (2) and Supplemental Material, Sec. 5 [33]]. The resulting $G_{\text{AFI}}(t)$, convoluted with our time resolution $\delta_{\text{exp}}(t)$, are shown in Fig. 3(a) for $d=3, 8, 18$, and 30 nm. To quantitatively assess the impact of d on the spin conductance, we fit a Gaussian function to the peak-shaped $G_{\text{AFI}}(t)$ (Supplemental Material, Sec. 6 [33]) and extract its amplitude [Fig. 3(b)], full width at half maximum (FWHM) [Fig. 3(c)] and time delay t_{pk} [Fig. 3(d)]. The amplitude decreases monotonically on a length scale of 10 nm [Fig. 3(b)]. In contrast, the FWHM increases monotonically for $d \geq 8$ nm, reaching 0.36 ps at 30 nm, but exhibits a local maximum at $d=3$ nm [Fig. 3(c)]. The peak time t_{pk} increases approximately linearly with d for $d<10$ nm and slightly sublinearly for $d>10$ nm [Fig. 3(d)].

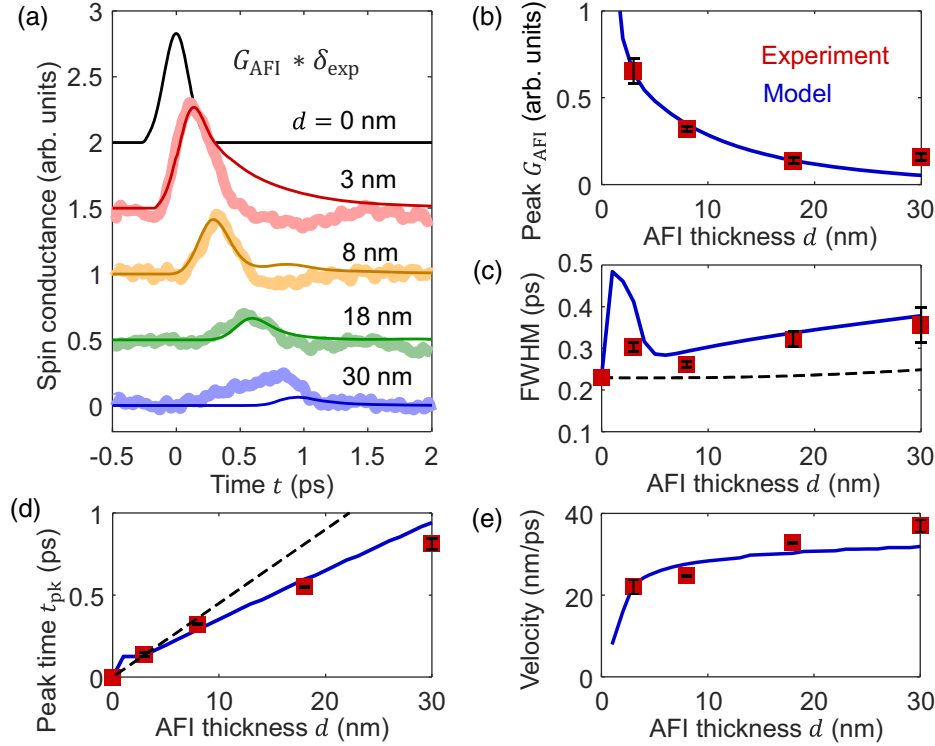


FIG. 3. (a) Time-domain spin conductance ($G_{\text{AFI}} * \delta_{\text{exp}}$)(t) of the NiO layer for various NiO thicknesses d . The convolution with $\delta_{\text{exp}}(t)$ accounts for the experimental time resolution. Marks and solid lines are measured and modeled data, respectively. (b) Extracted amplitude, (c) FWHM, (d) peak time t_{pk} , and (e) mean velocity d/t_{pk} of the peaklike ($G_{\text{AFI}} * \delta_{\text{exp}}$)(t) as a function of the NiO thickness d . The dashed line in (c) is the expected FWHM of coherent magnon transport. The dashed line in (d) is a linear extrapolation of the data points for $d < 10$ nm. The uncertainty bars in (b)–(e) represent the standard deviations of multiple measurements.

Consequently, the resulting signal velocity d/t_{pk} increases from 22 nm/ps at 3 nm to 37 nm/ps at 30 nm [Fig. 3(e)].

Interpretation—We emphasize that the velocities shown in Fig. 3(e) are consistent with the magnon group velocities of NiO at wavevectors exceeding 5% of the Brillouin-zone extent [see dashed lines in Fig. 4(a)]. This agreement is significant because, as a magnetic insulator, NiO permits spin transport primarily by magnons, tunneling, or by metallic imperfections like pinholes. For comparison, in insulating MgO, spin-polarized electron transport through pinholes or tunneling decays by more than 1 order of magnitude within just 1 nm [26], which is 2 orders of magnitude more rapid than that of G_{AFI} observed in NiO [Fig. 3(b)]. These results support the interpretation that the measured $G_{\text{AFI}}(t)$ [Fig. 3(a)] predominantly reflects spin transport by magnons and triggered by the Py spin voltage. A possible Seebeck-type contribution due to the pump-induced increase of the Py electron temperature can be ruled out (Appendix D).

Based on these findings, we propose the following scenario as the dominant mechanism behind our observations. Optical pumping transiently raises the electronic temperature by ΔT_e in the FM, while leaving the AFI unexcited. The ΔT_e results in a spin voltage μ_s (Fig. 1) [23], which injects spin angular momentum and, thus, magnons into the AFI through spin-transfer torque. Once the magnons

reach the AFI/HM interface, a fraction of the magnonic spin current enters the HM through spin pumping.

We can draw three key conclusions regarding magnon transport. First, if the transport was predominantly diffusive, one would expect the appearance of a long diffusive tail in the waveforms $G_{\text{AFI}}(t)$ [Fig. 3(a)] and a decrease in the average velocity d/t_{pk} , since diffusion implies $t_{\text{pk}} \propto d^2$ [Fig. 3(e)]. However, the absence of such features indicates that magnon transport is more ballisticlike. Second, specular reflections at the FM/AFI and AFI/HM interfaces are expected to occur, given the estimated reflection coefficients of about 0.6 and 0.8, respectively (Supplemental Material, Sec. 7 [33]). However, the lack of clearly discernible echoes in the experimental data [Fig. 3(a)] suggests that magnon reflections at the interfaces are likely diffuse rather than specular. Third, the transported magnons could, in principle, be either spatially coherent or incoherent. Coherent magnons propagate akin to plane waves normal to the FM/AFI interface. Given the nearly linear magnon dispersion of NiO within the frequency range covered [Fig. 4(a)], the magnon pulses would broaden by less than 10% over a 30 nm propagation distance in NiO [Fig. 3(c) and Supplemental Material, Sec. 8 [33]]. In contrast, the observed broadening exceeds 50% [Fig. 3(c)], implying

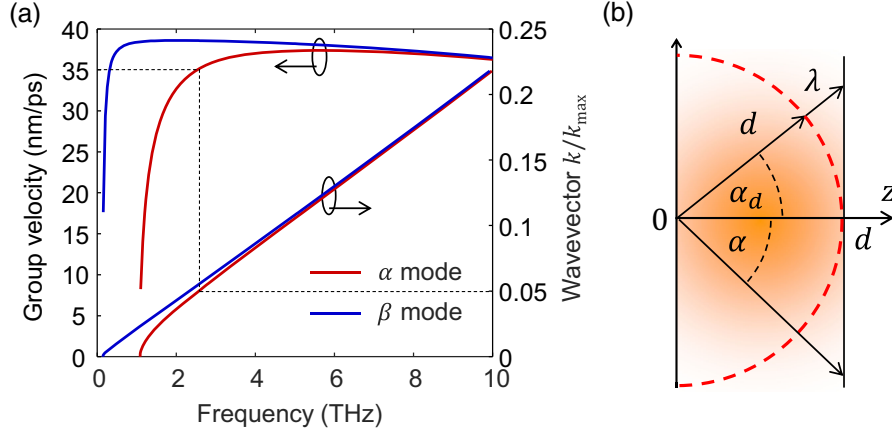


FIG. 4. Incoherent magnon propagation. (a) Magnon dispersion relation of NiO. Wavevector $k = |\mathbf{k}|$, normalized by the maximum wavevector $k_{\max}$, and group velocity $\partial\Omega_k/\partial k$ vs frequency $\Omega_k/2\pi$. (b) Schematic of the propagation of the angular momentum of an incoherent magnon pulse originating from the plane at $z = 0$. The angular momentum propagates spherically into the half-space $z > 0$ with group velocity v . Such spherical waves are launched throughout the laser-excited region in the x - y plane. Propagation directions $\alpha > \alpha_d$ do not contribute substantially to the angular momentum arriving at $z = d$ because their propagation lengths exceed that of the $\alpha = 0^\circ$ direction by at least the magnon relaxation length $\lambda = v\tau$.

that the measured $G_{\text{AFI}}(t)$ [Fig. 3(a)] are primarily carried by spatially incoherent magnons.

Model vs experiment—To quantitatively assess our interpretation, we develop a theoretical description of incoherent magnon transport based on the linearized Boltzmann transport equation (see Appendix E). By choosing a global magnon relaxation time of $\tau = 0.6$ ps, the calculated spin current $(G_{\text{AFI}} * \delta_{\text{exp}})(t)$ [Fig. 3(a), solid line] in response to a spin voltage $\delta_{\text{exp}}(t)$ agrees well with the measurement (marks) for all NiO thicknesses. An exception is the discrepancy for the largest thickness ($d = 30$ nm). We ascribe it to the increasing contribution of diffusive transport that is not accounted for in our model. Nevertheless, our calculations well capture the d dependence of the various parameters characterizing $(G_{\text{AFI}} * \delta_{\text{exp}})(t)$ [Figs. 3(b)–3(e)].

Discussion—Our model provides important insights into the ultrafast magnon transport in NiO, as detailed in Supplemental Material, Sec. 7 [33]. First, the amplitude decay of the magnon pulse $G_{\text{AFI}}(t)$ as a function of d [Fig. 3(b)] is dominated by the relaxation length $\lambda = v\tau$ of 20 nm, corresponding to the product of group velocity and relaxation time of the magnons. The extracted λ agrees excellently with the order of magnitude of relaxation lengths found in previous works on incoherent magnon transport in NiO thin films [10–12, 14–16]. However, the extracted τ is one order of magnitude shorter than the relaxation time of the uniform magnon mode at ≈ 1 THz in similar samples [28, 54]. We ascribe the faster relaxation to the fact that, in our experiment, the propagating incoherent magnons cover a much larger range of wavevectors and frequencies and have, thus, much more phase space for scattering.

Second, the FWHM of $G_{\text{AFI}}(t)$ increases sublinearly with d , because the mutually incoherent magnon wave packets can travel at all angles α relative to the FM/AFI-interface normal [Fig. 4(b)]. Magnons cover all velocities along the z axis, from

v ($\alpha = 0^\circ$) down to 0 ($\alpha = \pm 90^\circ$), which form the leading and trailing edges of $G_{\text{AFI}}(t)$ [55]. However, only magnons with $|\alpha| < \alpha_d$ contribute significantly, as those with $|\alpha| > \alpha_d$ must travel more than one relaxation length $\lambda = v\tau$ to reach the interface. Under this constraint, the FWHM scales approximately as $(d/\cos\alpha_d - d)/v$. As α_d decreases with increasing d , the FWHM exhibits a sublinear growth.

Third, the nonmonotonic behavior of FWHM as a function of d between 0 and 10 nm [Fig. 3(c)] arises from magnon reflection at the NiO/Pt interface and convolution with our finite time resolution $\delta_{\text{exp}}(t)$ [Fig. 3(a)]. The reflected magnon pulse temporally overlaps with the initial pulse, and this overlap is further blurred by the convolution, leading to an apparent broadening of the signal. The convolution-induced blurring also results in a nonmonotonic dependence of the delay time as predicted by the model [Fig. 3(d)]. Further, convolution in combination with a wider spread between leading and trailing edge may result in a shift of the peak position and, thus, increase of the propagation velocity with rising d [Fig. 3(e)].

Conclusion—We observe spintronic THz emission from Py|NiO|Pt stacks, where the thickness of NiO significantly influences the emitted THz waveforms. Our data can be consistently explained by ultrafast spin-voltage-driven spin transport through NiO. The spin current is predominantly carried by a ballistic flow of incoherent magnons at a high velocity of up to 40 nm/ps, which approximately equals the maximum magnon group velocity in NiO. The relaxation time of 0.6 ps permits ballistic propagation lengths of 20 nm, which is 1 order of magnitude larger than for THz spin-current pulses of conduction electrons in metals, such as Cu [25]. We even expect that the magnon relaxation length in NiO can be increased considerably by decreasing the defect density of the material. Indeed, the relaxation time of the uniform 1 THz magnon mode in NiO single

crystals [56] is about 1 order of magnitude longer than in NiO thin films [28,54]. From a methodological viewpoint, our Letter highlights the potential of spin-conductance spectroscopy of FM|AFI|HM stacks as a versatile platform to study THz magnon transport in any insulator with any magnetic order.

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Data availability—The data that support the findings of this article are openly available [57].

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End Matter

Appendix A: Experimental details—We study FM|AFI|HM|sub stacks, where Py(3 nm) serves as the FM, NiO as the AFI, and Pt(4 nm) as the HM on a MgO(500 μm) substrate (sub). An Al capping layer is applied for protection (Supplemental Material, Sec. 1 [33]). The NiO has a thickness d ranging from 0 to 30 nm. The samples are excited with linearly polarized femtosecond pump pulses from a Ti:sapphire oscillator (nominal duration 10 fs, energy 3 nJ, photon energy 1.6 eV, repetition rate 80 MHz). The emitted THz electromagnetic field is detected by electro-optic sampling in a GaP crystal (thickness 250 μm , bandwidth 1–7 THz) or ZnTe crystal (1 mm, 1–5 THz, yet higher sensitivity), yielding the electro-optic signal $S(t, M)$ [33]. All measurements are performed at room temperature in dry air. The FM magnetization M is controlled by an external magnetic field of ± 100 mT. We focus on the signal component

$$S(t) = \frac{S(t, +M) - S(t, -M)}{2}, \quad (\text{A1})$$

odd in M because it includes the spin transport indicated in Fig. 1. The even-in- M component of the THz signals is much weaker (Supplemental Material, Sec. 2 [33]) and attributed to off-resonant AFI spin pumping in the NiO/Pt regions [54,58].

Appendix B: Finite bandwidth of the response function $R(t)$ —Formally, $R(t)$ is the inverse Fourier transformation of $\tilde{R}(\omega)$. To restrict the analysis to the experimentally accessible bandwidth of 0–5 THz, a Norton-Beer window function $\tilde{\delta}_{\text{exp}}(\omega)$ is applied to $\tilde{R}(\omega)$ [Fig. 2(b)] [59]. The corresponding time-domain function $\delta_{\text{exp}}(t)$ is a unipolar pulse with a width of 240 fs [Fig. 2(c)]. The $\delta_{\text{exp}}(t)$ serves as an experimental approximation of the ideal $\delta(t)$ function within our time resolution [Fig. 2(c)] [60].

The filtered spectra $\tilde{R}(\omega)\tilde{\delta}_{\text{exp}}(\omega)$ [Fig. 2(b)] yield the time-domain signals $(R * \delta_{\text{exp}})(t)$ [Fig. 2(c)]. The convolution

$$(R * \delta_{\text{exp}})(t) = \int dt' R(t') \delta_{\text{exp}}(t - t') \quad (\text{B1})$$

smooths out any features of R finer than the width of δ_{exp} , while preserving its overall temporal shape. Details of the numerical implementation of Eq. (1) are provided in Supplemental Material, Sec. 3 [33].

If the THz-field detector of ZnTe is replaced by a more broadband GaP detector, the width of $\delta_{\text{exp}}(t)$ [Fig. 2(c)] reduces from 240 to 120 fs (Supplemental Material, Sec. 4 [33]).

Appendix C: Technical aspects of the extracted $(R * \delta_{\text{exp}})(t)$ —Here, we address three aspects of our response-function analysis [Eqs. (1) and (2)]. First, the extraction of $R(t)$ is essential because the two-peak structure of $(R * \delta_{\text{exp}})(t)$ [Fig. 2(c)] is barely discernible in the raw signals $S(t)$ [Fig. 2(a)]. Only when a THz detector with larger bandwidth is used, the two peaks appear in $S(t)$, but at the expense of less signal-to-noise ratio (Supplemental Material, Sec. 4 [33]). Second, even though the peak at $t = 0$ ps carries no information about spin transport through NiO, it is essential to precisely set the time-axis origin of $R(t)$.

Third, if the detection bandwidth was reduced by a factor of 2 or more, the two peaks of $R(t)$ would broaden and merge into a single peak. The maximum of the merged peak would still move toward later times $t > 0$ ps with increasing NiO thickness, but substantially less than the second peak. This effect can be perceived as an apparently shorter propagation time, potentially giving the impression of the superluminal magnon propagation suggested previously [28] (Supplemental Material, Sec. 4 [33]).

Appendix D: Possible Seebeck-type contribution—The total spin current is $j_s = j_s^{\Delta T_e} + j_s^{\mu_s}$, where the Seebeck-type contribution $j_s^{\Delta T_e}$ is driven by the pump-induced change ΔT_e in the FM electron temperature, and $j_s^{\mu_s}$ is driven by the FM spin voltage μ_s . Note that $\Delta T_e(t)$ has very different temporal dynamics than $\mu_s(t)$. In fact, within our time resolution, i.e., after convolution with $\delta_{\text{exp}}(t)$, $\mu_s(t)$ approximately scales with the time derivative of $\Delta T_e(t)$ [23]. Thus, we have $\tilde{\mu}_s(\omega) \propto i\omega \tilde{\Delta T}_e(\omega)$ in the frequency domain.

If the total j_s was dominated by the Seebeck contribution $j_s^{\Delta T_e}$, we would define and obtain the associated spin conductance in the frequency domain according to

$$\tilde{G}_{\text{AFI}}^{\Delta T_e}(\omega) = \frac{\tilde{j}_s^{\Delta T_e}(\omega)}{\tilde{\Delta T}_e(\omega)} = \frac{\tilde{j}_s(\omega)}{\tilde{\Delta T}_e(\omega)} \propto i\omega \frac{\tilde{j}_s(\omega)}{\tilde{\mu}_s(\omega)} = i\omega \tilde{G}_{\text{AFI}}^{\mu_s}(\omega). \quad (\text{D1})$$

Therefore, in the time domain, $G_{\text{AFI}}^{\Delta T_e}(t)$ would scale with the derivative $\partial_t G_{\text{AFI}}^{\mu_s}(t)$ of the spin conductance $G_{\text{AFI}} = G_{\text{AFI}}^{\mu_s}$ that was erroneously ascribed to a μ_s -driven process. Consequently, $G_{\text{AFI}}^{\Delta T_e}(t)$ would be proportional to the time derivative of the waveforms in Fig. 4(a) of the main text and, thus, assume a bipolar shape. This behavior is, however, unphysical: because the ΔT_e -driven spin-current density injected into the NiO is proportional to the unipolar $\Delta T_e(t) \propto \delta(t)$ [61], $G_{\text{AFI}}^{\Delta T_e}(t)$ should possess a unipolar shape also.

Appendix E: Brief description of the theoretical model—Our model relies on the linearized Boltzmann transport equation. Spin injection by the pump-induced spin voltage μ_s in the FM is modeled by spin-transfer torque exerted on the AFI. The generated incoherent magnons in the AFI propagate ballistically. Scattering is incorporated through a finite global relaxation time τ of the population of each magnon mode, while scattered magnons and, thus, diffusive transport are neglected. Reflections at the FM/AFI and AFI/HM interfaces are modeled as diffuse. Finally, spin transfer from the AFI

into the HM is described by spin pumping, leading to the spin-current density $j_s(t)$. The parameters of the simplified one-band magnon dispersion of NiO are taken from literature, and the spin-mixing conductance at the Py/NiO and NiO/Pt interfaces is estimated using the Sharvin limit and the value of YIG/Pt, respectively. A magnon relaxation time of $\tau = 0.6$ ps provides the best agreement with the experimental data. For details and a model of purely coherent magnon transport, we refer to Supplemental Material, Sec. 7 [33].