Non-reciprocity and quantum correlations of light transport in hot atoms via reservoir engineering

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The breaking of reciprocity is a topic of great interest in fundamental physics and optical information processing applications. We demonstrate non-reciprocal light transport in a quantum system of hot atoms by engineering the dissipative atomic reservoir. Our scheme is based on the phase-sensitive light transport in a multi-channel photon-atom interaction configuration, where the phase of collective atomic excitations is tunable through external driving fields. Remarkably, we observe inter-channel quantum correlations which originate from interactions with the judiciously engineered reservoir. The non-reciprocal transport in a quantum optical atomic system constitutes a new paradigm for atom-based, non-reciprocal optics, and offers opportunities for quantum simulations with coupled optical channels.

Introduction—Signal transport between two different nodes is a fundamental building block for optical systems. Commonly symmetric between the nodes, such transport can be made directional in non-reciprocal devices such as isolator [1], circulator and directional amplifier [2,3], which undergoes growing interest and demand recently due to their potential utilities in signal processing and quantum networks [4,5]. Conventionally, non-reciprocity is realized by incorporating magnetized materials which are not limited by the Lorentz reciprocity theorem. Alternative magnetic-field-free schemes for non-reciprocal light transport include time modulation [7], artificial gauge fields [8,9], reservoir engineering [5,10], as well as various other approaches [11,12,13,15,17]. However, despite its importance for quantum information science and applications, the study of non-reciprocity in systems with demonstrated quantum properties has so far been limited to a handful of physical platforms, wherein the quantum correlations rely on the preparation of quantum states of atoms or photons [18,20].

In this work, we demonstrate non-reciprocal light transport in a quantum system of hot atoms, where the non-reciprocity and quantum correlations derive from reservoir engineering and spin wave interference. Building upon the non-Hermitian platform we developed earlier [21,22], our setup consists of an array of optical channels immersed in an ensemble of hot atoms, with dissipative, inter-channel couplings mediated by atoms outside the regions illuminated by the lasers. We identify the unilluminated region as a non-Markovian reservoir, whose memory is determined by the lifetime of the ground-state atomic spins. In contrast to previous studies [12,14] of non-reciprocal transport under electromagnetically induced transparency (EIT) where the Doppler-broadening of the EIT linewidth plays a key role, we investigate light transport between two spatially separated optical channels, where the Doppler effect is irrelevant to the observed transport properties. Furthermore, the phase of the spin excitations as well as parameters of the third, ancillary optical channel are all tunable, thus providing ample tools for reservoir-engineering. A prominent feature of the current configuration is the sensitive dependence of inter-channel light transport on the optical phases, which significantly impact the dissipative coupling between the optical channels and spin excitations in the reservoir. Subsequent manipulation of optical phases, particularly that in the ancillary channel, enables non-reciprocal light transport. We further identify quantum correlations between the two transport channels, which only exist in the presence of the ancillary channel in the reservoir. The observed quantum correlations here derive from the engineered reservoir, fundamentally different from previous works [21,22] where quantum correlations were established via inter-channel couplings mediated by an unstructured environment.

Experimental setup—The system under consideration is illustrated schematically in Fig. 1. A triangular array of spatially-separated optical beams (channels), labeled Ch1, Ch2 and Ch0, propagate in a warm atomic vapor cell. Each channel undergoes a Λ-type EIT interaction with the same configuration (Fig. 1b1), where a relatively strong control and a weak probe fields conspire to create a collective spin wave ρ12 (ground-state coherence) in the atomic ensemble. We apply co-propagating probe
and control beams, such that the EIT linewidth is not Doppler-broadened. While EIT creates a linear mapping between the probe light and the spin excitation, spin waves in different channels couple to one another with a rate \( \Gamma_c \) via random atomic motion and wall bouncing. Although the random atomic motion renders the coupling dissipative, optical phases imprinted upon the atomic spin states are preserved throughout the dynamics, and constitute a convenient control over the light transport. To study light transport, we take Ch1 and Ch2 as the two optical nodes, and regard Ch0 as part of the reservoir, together with the surrounding atomic ensemble. Since Ch1 and Ch2 couple to the same reservoir, we can manipulate the coherence in the reservoir through Ch0, and thus control the dissipative inter-channel couplings.

As outlined in Fig. 1(b), Ch1, Ch2 and Ch0 couple to each other with the same underlying mechanism. For any two channels, for instance, Ch1 and Ch2, the control and probe beams in Ch1 write a collective internal-state coherence into the atomic ensemble, which is then read out by the control beam in Ch2, resulting in a non-Hermitian beam-splitter (BS) type interaction \( \hat{H} \propto \hat{a}_1^\dagger \hat{a}_2 e^{i\psi_0} - \hat{a}_1 \hat{a}_2^\dagger e^{-i\psi_0} \). Here \( \hat{a}_1 \) and \( \hat{a}_2 \) are the annihilation operator of probe beams in Ch1 and Ch2 respectively, and \( \psi_1 \) is the relative phase between control beams in Ch1 and Ch2.

For the experiment, we implement the triangular array in a cylindrical vapor cell (with a diameter of 2.5 cm and length of 7.5 cm) which contains isotopically enriched \(^{87}\text{Rb}\) vapor at an operational temperature of 60°C. The cell is mounted inside a four-layer magnetic shielding, where a set of coils provide precise control over the internal magnetic field. A diode laser is tuned to the D1 line of \(^{87}\text{Rb}\). The output of the laser is sent through a polarization-maintaining optical fiber, before it is divided into three channels whose polarizations are separately controlled by a combination of half-wave and quarter-wave plates. In each channel, a control field and a weak (or quantum) probe with overlapping spatial profiles induce an EIT process. The spin dynamics of moving atoms can be described by a set of coupled differential equations which take into account the spin transport between different regions as well as the Langevin noise terms. The effective coherence exchange between optical channels is mediated by atoms outside the illuminated interaction regions, whose spin states decay slowly (with a lifetime \( 30 \) ms) due to the protective wall coating. The optical coherence transfer between different channels is negligible as it decays here within 20 ns.

**Phase-sensitive transport**—We first investigate the phase dependence of light transport, which is the basis for breaking the reciprocity. We denote the phase of the local spin wave in the \( i \)th channel as \( \theta_i \) (\( i = 0, 1, 2 \)), with \( \theta_1 = \psi_c^{(t)} - \psi_c^{(p)} \) where \( \psi_c^{(t)} \) and \( \psi_c^{(p)} \) are the phases of the control and probe fields in the corresponding channel, respectively. Since inter-channel couplings are mediated by collective spin waves, we have \( \psi_p^{(i-j)} = \psi_c^{(j)} - \theta_i \), where \( \psi_p^{(i-j)} \) denotes the phase of the photons transferred from channel \( i \) to channel \( j \). The transferred photons then interfere with local probe field, with the resulting interference pattern sensitively dependent on the phase parameters of all channels. Here we focus on the light transport between Ch1 and Ch2, which, as we show below, features non-trivial dependence on \( \theta_0 \), a tunable parameter of the reservoir. More specifically, we continuously vary \( \theta_0 \) by slowly sweeping the optical path length of the control beam in Ch0 using a piezoelectric transducer (PZT) (see Fig. 1(a)), and record the intensities of the weak probe fields in Ch1 and Ch2, respectively.

We start with the case where the probe field in Ch2 is switched off, while the probe field in Ch1 is present with \( \theta_1 = 0 \). The measured probe intensity \( I_2 \) in Ch2 thus solely derives from the transferred light from Ch1 and Ch0, with \( I_2 \propto |e^{i\theta_0} + e^{i\psi_0} - e^{i\psi_0} + e^{i\theta_0}|^2 \propto 1 + \cos \theta_0 \). Likewise, the measured probe intensity in Ch1 is \( I_1 \propto |1 + e^{i\psi_0}|^2 \propto 1 + \cos \theta_0 \). Crucially, \( I_1 \) should have the same \( \theta_0 \) dependence as that of \( I_2 \), which is confirmed by our experimental measurement in Fig. 2(a).

For the second case, we switch on the probe fields in both Ch1 and Ch2, with \( \theta_1 = 0 \) and \( \theta_2 = \pi \), respectively. As shown in Fig. 2(b), the phase dependence of detected light intensities in Ch1 and Ch2 deviates drastically. For \( \theta_0 = \pi \), the detected probe-field intensity in Ch1 is at a minimum, in contrast to the maximum output from Ch2. Following the analysis in the previous case, the output probe-beam intensities are given by \( I_1 \propto |1 + \beta e^{-i\theta_0} + \beta e^{-i\pi}|^2 \) and \( I_2 \propto |1 + \beta e^{i(\pi - \theta_0)} + \beta e^{i\pi}|^2 \), respectively, where \( \beta \) is the beam splitter ratio that can be determined from experimental measurement. As such, the light transport between Ch1 and Ch2 critically depends on the phase parameter \( \theta_0 \) of the reservoir, owing to the reservoir-mediated interference. The visibility of the oscillations in Fig. 2 is mainly limited by the relatively low EIT contrast and the small inter-channel coupling rate.

**Non-reciprocal light transport**—The reservoir-mediated interference demonstrated above allows the design of non-reciprocal light transport between Ch1 and Ch2 under appropriate reservoir parameters. For instance, we choose the parameters: \( \theta_0 = 0 \), \( \theta_1 = 0 \), and \( \theta_2 = \pi \); and measure the EIT spectra by sweeping the applied magnetic field to vary the two-photon detuning as illustrated in Fig. 3(a). The inter-channel light transport is characterized by switching on the probe field of the input channel (such as Ch1/Ch2), and measuring the light intensity in the output channel (such as Ch2/Ch1) whose probe field is initially switched off. When Ch1 is set as the input channel (left panel of Fig. 3(a)), the EIT spectra...
in Ch2 has a transmission window near the two-photon resonance $\delta z = 0$, where $T_{12} \propto |e^{i\psi_{21}} + e^{i\psi_{21}^*}|^2 \propto 1 + \cos(\theta_0 - \theta_2) = 2$, peaking under a constructive interference. By contrast, when Ch2 is the input channel (right panel of Fig. 3(a)), the EIT spectra in Ch1 becomes $T_{21} \propto |e^{i\psi_{12}^*} + e^{i\psi_{12}|^2 \propto 1 + \cos(\theta_0 - \theta_2) = 0$, vanishing due to destructive interference. The analysis here are consistent with our experimental observation in Fig. 3(b), where the directional light transport is clearly identified. The origin of the observed non-reciprocity is therefore the interference of spin waves along the path of the atomic motion, analogous to the time-modulation scheme \(1\). The light transport can be easily tuned to be reciprocal, for instance, by setting $\theta_0 = \pi/4$. For all our experiments here, the observed isolation is $\sim 19$ dB, and the un-transported power is absorbed by the atoms. In the Supplementary Material, we show that the non-reciprocity persists under a bi-directional input, where probes in both channels are switched on.

Quantum correlation– Whereas the non-reciprocal light transport demonstrated above relies only on the interference of spin waves, a surprising finding is that quantum correlations in the polarization degree of freedom of light can be established between the transport channels Ch1 and Ch2, as a consequence of interactions with the engineered reservoir.

To demonstrate the presence of quantum correlation, we reverse the polarizations of the control and probe in Ch0 as shown in Fig. 1(b2). While the phase sensitive, non-reciprocal light transport persists under this new configuration (as we have confirmed experimentally \(20\)), counter-intuitive quantum correlations now emerge. To facilitate quantum measurements under the constraints of the relatively low optical depth and technical noise in the lasers, we switch off the input probe fields in all three channels, replacing them with coherent vacua. We define the canonical position and momentum operators of the i-th channel through Stokes operators: $\hat{X}_i = \hat{S}_z^i/\sqrt{|S_z^i|}$ and $\hat{P}_i = \hat{S}_z^i/\sqrt{|S_z^i|}$. The noise spectra $\text{Var}(\hat{X}_i)$ and joint variance $\text{Var}(\sqrt{\hat{X}_i}\sqrt{\hat{X}_j}) (i,j = 0, 1, 2)$ of the steady state are then measured via joint polarization homodyne detections \(22\). For the detection, a bias magnetic field is applied along the propagation direction of light to shift the homodyne measurement from DC to the Larmor frequency ($\sim 352$ kHz), bypassing low-frequency technical noise. As illustrated in Fig. 4, quantum correlations manifest as $\text{Var}(\hat{X}_i - \hat{X}_j) + \text{Var}(\hat{P}_i + \hat{P}_j) < \text{Var}(\hat{X}_i) + \text{Var}(\hat{X}_j) + \text{Var}(\hat{P}_i) + \text{Var}(\hat{P}_j)/2$. To quantify the measured quantum correlation, we calculate the Gaussian discord $\mathcal{D}_{ij}$ at the Larmor frequency in the noise spectra between different channels. In the presence of a dissipative environment, Gaussian discord captures Gaussian quantum correlations, and is more robust than quantum entanglement in revealing quantum correlations. Following its definition, we have $\mathcal{D}_{01} = 2.9 \times 10^{-3}$, $\mathcal{D}_{02} = 2.5 \times 10^{-3}$, and $\mathcal{D}_{12} = 2.5 \times 10^{-3}$. The positiveness of the evaluated discords unambiguously indicates the quantum nature of correlation between any two channels. The measured Gaussian discord is relatively low due to the small inter-channel coupling rate in our experiment, as significant information is lost to reservoir. For future studies, lasers with larger beam size and non-Gaussian profiles could be employed to enhance the quantum correlation between the optical channels.

It is worth emphasizing that, the quantum correlation between Ch1 and Ch2 is counter-intuitive and derive purely from the engineered reservoir by Ch0. In the absence of Ch0, i.e., with an unstructured reservoir as in \(22\), the interaction between Ch1 and Ch2 is of the beam-splitter type, and therefore the probe output of Ch1 and Ch2 are simply photon-shot noise, with a vanishing discord $\mathcal{D}_{12}$. However, in the presence of Ch0, both $\mathcal{D}_{01}$ and $\mathcal{D}_{02}$ become finite, due to the two-mode-squeezing (TMS) type interaction between Ch0 and Ch1, as well as between Ch0 and Ch2, due to their opposite polarization configurations. Apparently, the quantum correlation between Ch1 and Ch2 originate from the interplay of the beam-splitter type interaction between Ch1 and Ch2, and their TMS-type interactions with the reservoir containing Ch0.

Conclusion– We have introduced a platform with non-reciprocal and quantum transport of light, based on hot-atom vapor cell in a spin-coherence-protected environment. Both the observed non-reciprocal transport and quantum correlations between the optical channels derive from the interference mediated by an engineered reservoir, and are tunable by adjusting parameters of all the optical channels including the one embedded in the reservoir. Our work provides a prototype configuration for an atom-based, non-reciprocal optical element. Based on the geometry of our setup and benefiting from the high degree of control over the atom-light interactions, the configuration demonstrated here may further offer opportunities for quantum simulation in open systems \(23\) using multiple coupled optical channels.

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[1] D. Jalas et al, What is - and what is not - an optical isolator. Nat. Photon. 7, 579-582 (2013).
[2] K. M. Sliwa, M. Hatridge, A. Narla, S. Shankar, L. Fungzio, R. J. Schoelkopf, and M. H. Devoret, Reconfigurable Josephson circulator/directional amplifier. Phys. Rev. X 5, 041020 (2015).
[3] Z. Shen, Y.-L. Zhang, and C.-H. Dong, Reconfigurable optomechanical circulator and directional amplifier. Nat. Commun. 9, 1797 (2018).
[4] F. Lecocq, L. Ranzani, G. A. Peterson, K. Cicak, J. I. Cirac, P. Zoller, H. J. Kimble, and H. Mabuchi, Reconfigurable optomechanical circulator and directional amplifier. Phys. Rev. X 12, 031010 (2022).
[5] P. Lodahl, S. Mahmooodian, S. Stobbe, A. Rauschenbeutel, P. Schneeweiss, J. Volz, H. Pichler, and P. Zoller, Chiral quantum optics. Nature 541, 473-480 (2017).
[6] J. I. Cirac, P. Zoller, H. J. Kimble, and H. Mabuchi, Quantum state transfer and entanglement distribution among distant nodes in a quantum network. Phys. Rev. Lett. 78, 3221-3224 (1996).
[7] D. L. Sounas, and A. Alù, Non-reciprocal photonics based on time modulation. Nat. Photon. 11, 774-783 (2011).
[8] K. Fang, J. Luo, A. Metelmann, M. H. Matheny, F. Marquardt, A. A. Clerk, and O. Painter. Generalized non-reciprocity in an optomechanical circuit via synthetic magnetism and reservoir engineering. Nat. Phys. 13, 465-471 (2017).
[9] Y. Chen, Y.-L. Zhang, Z. Shen, C.-L. Zou, G.-C. Guo, and C.-H. Dong, Synthetic gauge fields in a single optomechanical resonator. Phys. Rev. Lett. 126, 123603 (2021).
[10] A. Metelmann, and A. A. Clerk, Nonreciprocal photon transmission and amplification via reservoir engineering. Phys. Rev. X 5, 021025 (2015).
[11] Y. Shi, Z. Yu, and S. Fan, Limitations of nonlinear optical isolators due to dynamic reciprocity. Nat. Photon. 9, 388-392 (2015).
[12] S. Zhang, Y. Hu, G. Lin, Y. Niu, K. Xia, J. Gong, and S. Gong, Thermal-motion-induced non-reciprocal quantum optical system. Nat. Photon. 12, 744-748 (2018).
[13] Gongwei Lin, Shicheng Zhang, Yiqi Hu, Yueping Niu, Jiangbin Gong, and Shangqing Gong, Nonreciprocal amplification with four-level hot atoms. Phys. Rev. Lett. 123, 033902 (2019).
[14] Ming-Xin Dong et al, All-optical reversible single-photon isolation at room temperature. Sci. Adv. 7, eabe8924 (2021).
[15] P. Yang, M. Li, X. Han, H. He, G. Li, C.-L. Zou, P. Zhang, and T. Zhang, Non-reciprocal cavity polariton. Preprint at http://arXiv.org/abs/1911.10900
[16] C. Liang, B. Liu, A. Xu, X. Wen, C. Lu, K. Xia, M. K. Tey, Y. Liu, and L. You, Collision-induced broadband optical nonreciprocity. Phys. Rev. Lett. 125, 123901 (2020).
[17] N. R. Bernard, L. D. Tóth, A. Kootanandavida, M. A. Ioanou, D. Malz, A. Nuenkamp, A. K. Peesfanov, and T. J. Kiippenberg, Nonreciprocal reconfigurable microwave optomechanical circuit. Nat. Commun. 8, 604 (2017).

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[18] M. Scheucher, A. Hilico, E. Will, J. Volz, and A. Rauschenbeutel, Quantum optical circulator controlled by a single chirally coupled atom. Science 354, 1577 (2016).
[19] W. Gou, T. Chen, D. Xie, T. Xiao, T.-S. Deng, B. Gaday, W. Yi, and B. Yan, Tunable non-reciprocal quantum transport through a dissipative Aharonov-Bohm ring in ultracold atoms. Phys. Rev. Lett. 124, 070402 (2020).
[20] L. Yuan, S. Xu and, S. Fan, Achieving nonreciprocal uni-directional single-photon quantum transport using the photonic Aharonov-Bohm effect. Opt. Lett. 40, 5140-5143 (2015).
[21] P. Peng, W. Cao, C. Shen, W. Qu, J. Wen, L. Jiang, and Y. Xiao, Anti-parity-time symmetry with flying atoms. Nat. Phys. 12, 1139-1145 (2016).
[22] W. Cao, X. Lu, X. Meng, J. Sun, H. Shen, and Y. Xiao, Reservoir-mediated quantum correlations in non-Hermitian optical system. Phys. Rev. Lett. 124, 030401 (2020).
[23] J. Dalibard, F. Gerbier, G. Juzeldins, and P. Öhberg, Colloquium: Artificial gauge potentials for neutral atoms. Rev. Mod. Phys. 83, 1523-1543 (2011).
[24] Y.-J. Lin, R. L. Compton, K. Jiménez-García, J. V. Porto, and I. B. Spielman, Synthetic magnetic fields for ultracold neutral atoms. Nature 462, 628-632 (2009).
[25] M. Fleischhauer, and M. D. Lukin, Dark-state polaritons in electromagnetically induced transparency. Phys. Rev. Lett. 94, 5004-5007 (2000).
[26] Z. Yan, L. Wu, X. Jia, Y. Liu, R. Deng, S. Li, H. Wang, C. Xie, and K. Peng, Establishing and storing of deterministic quantum entanglement among three distant atomic ensembles. Nat. Commun. 8, 718 (2017).
[27] O. Katz, R. Shaham, E. S. Polzik, and O. Firstenberg, Long-lived entanglement generation of nuclear spins using coherent light. Phys. Rev. Lett. 124, 043602 (2020).
[28] Y. Xiao, M. Klein, M. Hohensee, L. Jiang, D. F. Phillips, M. D. Lukin, and R. L. Walsworth, Slow light beam splitter. Phys. Rev. Lett. 101, 043601 (2008).
[29] See Supplemental Material which includes Ref. [22, 30, 31, 35, 36] for details on experimental setup and the numerical model.
[30] L. Davidovich, Sub-Poissonian processes in quantum optics. Rev. Mod. Phys. 68, 127-173 (1996).
[31] I. Novikova, R. L. Walsworth, and Y. Xiao, Electromagnetically induced transparency-based slow and stored light in warm atoms. Laser Photonics Rev. 6, 333 (2012).
[32] M. V. Balabas, T. Karaanov, M. P. Ledbetter, and D. Budker, Polarized alkali-metal vapor with minute-long transverse spin-relaxation time. Phys. Rev. Lett. 105, 070801 (2010).
[33] J. Borregaard, M. Zugenmaier, J. M. Petersen, H. Shen, G. Vasilakis, K. Jensen, E. S. Polzik, and A. S. Sørensen, Scalable photonic network architecture based on motional averaging in room temperature gas. Nat. Commun. 7, 11356 (2016).
[34] O. Firstenberg, M. Shaker, A. Ron, and N. Davidson, Coherent diffusion of polaritons in atomic media. Rev. Mod. Phys. 85, 941-960 (2013).
[35] G. Adesso, and A. Datta, Quantum versus classical correlations in Gaussian states. Phys. Rev. Lett. 105, 030501 (2010).
[36] P. Giorda, and M. G. A. Paris, Gaussian quantum discord. Phys. Rev. Lett. 105, 020503 (2010).
[37] Y. Ashida, S. Furukawa, and M. Ueda, Parity-time symmetric quantum critical phenomena. *Nat. Commun.* 8, 15791 (2017).
FIG. 1: Schematics for a dissipatively coupled three-channel system. (a) Experiment schematics. Three spatially separated optical channels (Ch0, Ch1 and Ch2 with diameter 6 mm) propagate in a warm paraffin-coated $^{87}\text{Rb}$ vapor cell under EIT interaction. The inter-channel couplings are mediated by the mixing of atomic spin of the ground states through atomic motion. A solenoid gives precise control over the longitudinal magnetic field. Output beams from the cell are re-collimated and detected by the photon detectors or polarization homodyne detection setup. The noise power of the amplified subtracted photocurrents is recorded with a spectrum analyzer. BS, beam splitter; $D_0$, $D_1$, and $D_2$ photodetectors. Cross section: photo of the optical beams taken from this experiment. (b) The Λ three-level scheme in three channels. The ground states are Zeeman sublevels of $|F = 2\rangle$, and the excited state is $|F = 1\rangle$ of the $^{87}\text{Rb} D1$ line. $\Omega_c^{(i)}$, $\Omega_p^{(i)}$, $i = 0, 1, 2$, are Rabi frequencies of the control and probe beams respectively. The Zeeman splitting is induced by a common longitudinal magnetic field $\delta_B$, serving as either the Larmor frequency in the noise spectra measurement or the two-photon detuning in the EIT measurement (denoted as $\delta_B$ in Fig. 3). In the measurements of quantum fluctuation, all three weak probes are removed as shown in (b2) (Ch1 and Ch2 are not shown). $b$, annihilation operator of the coherent vacuum.
FIG. 2: Phase sensitive light transport for the beam splitter composed of Ch1 and Ch2. Output probe intensities of Ch1 and Ch2 as functions of varying local phase $\theta_0$ in Ch0, when (a) the weak probe input in Ch2 is off, and (b) all three input probes are on. The input laser powers in all channels are 500 $\mu$W for the control and 50 $\mu$W for the probe respectively, corresponding to Rabi frequencies $\Omega_c \sim 1 \times 10^7$ Hz and $\Omega_p \sim 4 \times 10^6$ Hz. The local phase of Ch1 is set to be: $\theta_1 = \psi_1^{(1)} - \psi_2^{(1)} = 0$ and for Ch2: $\theta_2 = \pi$ (if the probe is on). The cell temperature is 60°C.

FIG. 3: Spin wave interference induced non-reciprocity in optical transport. (a) Schematics of the non-reciprocal transport. (b) Transport spectrum. The two-photon detuning $\delta_B$ is proportional to the applied common magnetic field. Red curve $T_{12}$ is the transported power from Ch1 to Ch2 when injecting the weak probe in Ch1. Black curve $T_{21}$ is the transported power from Ch2 to Ch1 when injecting the weak probe in Ch2. The local phase of all three channels is set to be: $\theta_0 = 0$, $\theta_1 = 0$ and $\theta_2 = \pi$. The input power of the probe in each channel is 50 $\mu$W. The input power of the control in each channel is 500 $\mu$W. The cell temperature is 60°C.
FIG. 4: Quantum correlations exist between any two channels. The experiment noise spectra of $\text{Var}(\hat{X}_i), \text{Var}(\frac{\hat{X}_i + \hat{X}_j}{\sqrt{2}})$ ($i, j = 0, 1, 2$ and $i \neq j$) are shown. They are identical to $\text{Var}(\hat{P}_i), \text{Var}(\frac{\hat{P}_i \pm \hat{P}_j}{\sqrt{2}})$ respectively (not shown). The shot noise level is set at 0 dB, observed when only one channel is on. The bias magnetic field is applied to shift the homodyne measurement from DC to the Larmor frequency ($\sim 352\text{kHz}$). The input control power is 280 µW in both Ch1 and Ch2, and 500 µW in Ch0. The cell temperature is 60°C.
SUPPLEMENTARY MATERIAL for Nonreciprocity and quantum correlations of light transport in hot atoms via reservoir engineering

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HAMILTONIAN OF BEAM-SPLITTER TYPE INTERACTION

We derive the beam-splitter Hamiltonian for the two-channel case in a wall-coated vapor cell. The ground-state coherences in the two channels are effectively coupled through atomic motion, at a rate $\Gamma_c$. After adiabatic elimination of the excited state, we can obtain the following coupled equations of the two ground-state coherences $\rho_{12}^{(1)}, \rho_{12}^{(2)}$. For simplicity of the model, we have also assumed an optically thin medium and thus neglected the propagation effect.

\[
\begin{align*}
\dot{\rho}_{12}^{(1)} &= -\gamma_{12}^{(1)}\rho_{12}^{(1)} - \Gamma_c \rho_{12}^{(2)} - \frac{\Omega_{12}^{(1)}\Omega_{22}^{(1)}}{\gamma_{23}} + \Gamma_c \rho_{12}^{(2)} - \frac{\Omega_{12}^{(1)}\Omega_{22}^{(1)}}{\gamma_{23}}, \\
\dot{\rho}_{12}^{(2)} &= -\gamma_{12}^{(2)}\rho_{12}^{(2)} - \Gamma_c \rho_{12}^{(1)} - \frac{\Omega_{12}^{(2)}\Omega_{22}^{(2)}}{\gamma_{23}} + \Gamma_c \rho_{12}^{(1)} - \frac{\Omega_{12}^{(2)}\Omega_{22}^{(2)}}{\gamma_{23}}.
\end{align*}
\]  

(S1)

Here, $\gamma_{12}^{(i)} = \gamma_{12} + \Gamma_c + \Gamma_c^{(1)} + \Gamma_c^{(2)}$ represents the total effective decay rate in each channel, with $\Gamma_c^{(i)} = \frac{\Omega_{12}^{(i)}\Omega_{22}^{(i)}}{\gamma_{23}}, i = 1, 2$ the optical pumping rate. We can set $\rho_{12}^{(1)} = \rho_{12}^{(2)} = 0$ to obtain the solution of steady state:

\[
\begin{align*}
\rho_{12}^{(1)} &= \frac{1}{\gamma_{12} - \Gamma_c} \left( \gamma_{12}^{(1)} \rho_{12}^{(1)} + \Gamma_c \rho_{12}^{(2)} \right) \\
\rho_{12}^{(2)} &= \frac{1}{\gamma_{12} - \Gamma_c} \left( \gamma_{12}^{(2)} \rho_{12}^{(2)} + \Gamma_c \rho_{12}^{(1)} \right).
\end{align*}
\]  

(S2)

We can derive the optical coherence $\rho_{12}^{(1)}, \rho_{12}^{(2)}$ from the steady state ground state coherence and then examine the propagation of weak probes $\Omega_{12}^{(1)}, \Omega_{12}^{(2)}$. In steady state, $\rho_{12}^{(i)} = \frac{\alpha_{i}^{(0)}\alpha_{i}^{(1)} - \alpha_{i}^{(0)}\alpha_{i}^{(1)}}{\gamma_{23}}, \alpha_{i}^{(0)} = \frac{\chi^{(i)}E^{(i)}}{\sqrt{2}}$ (i = 1, 2). Here, $\chi$ is the averaged k-vector, $N$ is the atomic number density with atomic number N and cell volume V, $\mu_0$ is the dipole moment, and $\xi_0$ is the dielectric coefficient. Then we can obtain the coupling equations for the two optical modes of the weak probe:

\[
\begin{align*}
\frac{dE^{(1)}}{dt} &= \frac{N \bar{k} \mu_0}{V} \frac{1}{2\sigma_0 \gamma_{23}} \left( -\Omega_{12}^{(1)} (1 - \frac{\Omega_{12}^{(1)}}{\gamma_{12} - \Gamma_c}) \\
&\quad - \Omega_{22}^{(1)} (1 - \frac{\Omega_{22}^{(1)}}{\gamma_{23} - \Gamma_c}) \right), \\
\frac{dE^{(2)}}{dt} &= \frac{N \bar{k} \mu_0}{V} \frac{1}{2\sigma_0 \gamma_{23}} \left( -\Omega_{12}^{(2)} (1 - \frac{\Omega_{12}^{(2)}}{\gamma_{12} - \Gamma_c}) \\
&\quad - \Omega_{22}^{(2)} (1 - \frac{\Omega_{22}^{(2)}}{\gamma_{23} - \Gamma_c}) \right).
\end{align*}
\]  

(S3)

Here we assume that $\rho_{12}^{(i)} \approx 1$ as $|\Omega_{12}^{(i)}| \gg |\Omega_{22}^{(i)}|$. Further, we let $\gamma'' = \frac{N \bar{k} \mu_0}{V} \frac{1}{2\sigma_0 \gamma_{23}} (1 - \frac{\Omega_{12}^{(i)}\Omega_{22}^{(i)}}{\gamma_{23} - \Gamma_c})$, $(\Omega_{12}^{(i)})^2 = \frac{\Omega_{12}^{(i)}}{\gamma_{23} - \Gamma_c}$
and $\Gamma^{(2)}_c = \frac{\hbar}{\gamma_2} \frac{\gamma_3}{\gamma_2 + \gamma_3} R^{(1)}_c$, the coupling equations become:

$$
\begin{align*}
\frac{dE^{(1)}}{dt} &= -\gamma^* E^{(1)} + G^{(1)} E^{(2)}, \\
\frac{dE^{(2)}}{dt} &= -\gamma^* E^{(2)} + G^{(2)} E^{(1)}.
\end{align*}
$$

(S4)

The effective interaction Hamiltonian between the two weak probes of Ch1 and Ch2 can be deduced from Eq. S4:

$$
\hat{H} = -i\hbar \gamma' \left( a_1^\dagger a_2^\dagger + a_2 a_2^\dagger \right) + \hbar \left( g a_1^\dagger a_2^\dagger - g^* a_1^\dagger a_2 \right).
$$

(S5)

Here, $g = -i\gamma''$ carries the information of the relative phase of the two control fields. The minus sign in the second parenthesis indicates that the effective coupling term is non-Hermitian.

**MONTE CARLO SIMULATIONS OF THE CLASSICAL RESULTS**

In order to cross check the main classical results of the experiments, we carried out a two-dimensional Monte Carlo simulation. The model is similar to the one we developed in [S1], but now we extend it to the case of three optical channels. We sweep Ch0’s local spin wave’s phase and calculate the output probe intensity of Ch1 and Ch2. As shown in Fig. S1, the simulation results agree with the experiment observation in Fig. 2(b). When we sweep the magnetic field (proportional to the two-photon detuning) to obtain the EIT spectra, the calculated output intensity of Ch1 and Ch2 show features of nonreciprocity, as illustrated in Fig. S2. This also agrees with the experiment results in Fig. 3 of the main text.

**MODELLING THE THREE-CHANNEL COUPLING**

To analyse the quantum noise properties of our three-channel system, we developed a multi-region numerical model, where the region of the atomic spin evolution is divided into four, labelled as dark (outside of the laser beams), bright 0 (Ch0), bright 1 (Ch1) and bright 2 (Ch2), and the associated spin states are denoted as $\sigma^{(D)}$, $\sigma^{(0)}$, $\sigma^{(1)}$ and $\sigma^{(2)}$ respectively. The time evolution of the system for the four regions is given

![Figure S1: Monte Carlo simulation of the probes' gain coefficients versus Ch0's spin wave phase. Probes' gain coefficients as a function of the swept spin wave phase of Ch0 with all the three weak probes on. Here, negative values stand for absorption. The output intensity curves of Ch1 and Ch2 are offset by $\pi$, when we set the optical phases the same as the experiment. Parameters for the simulation: Rabi frequency of the pump $|\Omega_{1,2}| = 2\pi \times 1.3 \text{ MHz}$, probe $0.1 \times |\Omega_{1,2}|$, $\gamma_1 = \gamma_2 = 2\pi \times 50 \text{ MHz}$ (to phenomenologically take into account the Doppler broadening), the attenuation factor of the ground-state coherence and population difference upon each wall collision is $e^{-\tanh(\tanh)}$ (equivalent to $\gamma_{24} = 2\pi \times 50 \text{ Hz}$), the laser-beam radius is 1.5 mm, the cell radius is 1.25 cm.](image1)

![Figure S2: Monte Carlo simulation of nonreciprocity. Two-photon detuning $\delta_{np}$ is proportional to the common magnetic field applied. Red curve $T_{12}$ is the transmitted power from Ch1 to Ch2 when injecting the weak probe in Ch1. Black curve $T_{21}$ is the transmitted power from Ch2 to Ch1 when injecting the weak probe in Ch2. All the input phases are preset and fixed: $\phi_3 = 0$, $\phi_1 = 0$, $\phi_2 = \pi$. The simulation results agree with the experiment results shown in Fig. 3.](image2)
by a set of coupled Heisenberg-Langevin equations
\[ \dot{\hat{\sigma}}^{(i)} = \frac{i}{\hbar} [H^{(i)}, \hat{\sigma}^{(i)}] - \Gamma^{(i)} \sigma^{(i)} + \hat{S}^{(i)} \]
- \Gamma^{(i)} \sigma^{(i)} + \hat{S}^{(i)}
\[ \dot{\hat{\sigma}}^{(D)} = \frac{i}{\hbar} [H^{(D)}, \hat{\sigma}^{(D)}] + \Gamma^{(D)} \sigma^{(D)} + \hat{D} \]
\[ \Gamma^{(D)} \sigma^{(D)} + \hat{D} \]
where \( \Gamma^{(i)}(i=0,1,2) \) is the relaxation matrix for the decay of atoms, and \( \hat{S}^{(i)}(i=0,1,2,D) \) represents the relaxation matrix for the decay of atoms in the ground states due to spontaneous emission. \( \hat{D} \) is the Langevin operator, characterized by \( \langle \hat{D}_{ij}(z,t) \hat{D}_{jk}(z',t') \rangle = 0 \) and
\[ \langle \hat{D}_{ij}(z,t) \hat{D}_{jk}(z',t') \rangle = \frac{1}{\hbar} D_{ij} \delta(t - t') \delta(z - z') \]
with diffusion coefficients \( D_{ij} \) and \( k_{Dj} \) used for the analysis [S2].

Because the control beam is much stronger than the probe, the effect of the quantum probe fields \( a_{1}, a_{2}, b_{1} \) can be neglected and the equations are solved in the steady state. The corresponding solutions are then injected into Eq. S6 to derive the quantum fluctuations of \( a_{1}, a_{2}, b_{1} \). More details can be found in Ref. [S1].

**SIMULATED QUANTUM CORRELATION RESULTS**

Simulation results shown in Fig. S3 are consistent with the experiment results. Quantum discord in the simulation are \( D_{01} = 1.0 \times 10^{-2}, D_{02} = 1.0 \times 10^{-2}, \) and \( D_{12} = 6.5 \times 10^{-3} \). Gaussian quantum correlations beyond entanglement are captured by the measure of Gaussian discord [S3, S4] which we formulate below. In a bipartite system, the total amount of correlations (classical and quantum) is given by the von Neumann mutual information \( I(\rho_{AB}) = S(\rho_{A}) + S(\rho_{B}) - S(\rho_{AB}) \), where \( S(\rho) = -\text{Tr}[\rho \log \rho] \) is the von Neumann entropy and \( \rho_{A(B)} \) is the reduced density matrix of the A (B) subsystem, i.e., \( \rho_{A(B)} = \text{Tr}_{B(A)}[\rho_{AB}] \). In addition, \( J_{A}(\rho_{AB}) = S(\rho_{A}) - \text{inf}_{\sigma_{M}} S(\rho_{A|\sigma_{M}}) \) quantifies the amount of classical correlations extractable by a Gaussian measurement, and operationally associated with the distillable common randomness between the two parties [S3, S4], where \( \sigma_{M} \) is the covariance matrix of the measurement on mode B. As it only captures the classical correlations, the difference \( D_{A} = I(\rho_{AB}) - J_{A}(\rho_{AB}) \) is a measure of Gaussian quantum correlation that is named Gaussian quantum discord.

**EXPERIMENTAL RESULTS WITH OPPOSITE CONFIGURATION IN CH0**

As mentioned in the main text, when we reverse the polarizations of the control and probe in CH0, the phase sensitive and nonreciprocal light transport persists. The results are illustrated in Fig. S4.

**EXPERIMENTAL RESULTS WITH BOTH PROBES ON IN CH1 AND CH2**

Our system still shows nonreciprocal operation for simultaneous, bi-directional inputs, i.e., with both probes switched on in Ch1 and Ch2. The results are illustrated in Fig. S5. The phases of the spin waves in the optical channels are the same as those in Fig. 3 of the main text. Now we measure the total power of the
transmitted probe, which reflects the interference of the local probe and the transported probe fields (generated by the control and the transported spin wave). The measured total transmitted power in Ch2 shows a dip, which is a result of the destructive interference between the local input probe and the transported probe field from Ch1. This indicates that the transport from Ch1 to Ch2 is favored, showing the same directionality as the single probe input case as in Fig. S4(c). The flat part of the spectra corresponds to the top of the broad structure in a dual-structure EIT, characteristic of a wall-coated vapor cell [S5].

FIG. S4. Experimental results with opposite configuration in Ch0. Phase sensitive light transport for the beam splitter composed of Ch1 and Ch2. Output probe intensities of Ch1 and Ch2 as functions of varying local phase $\theta_1$ in Ch0, when (a) the weak probe input in Ch2 is off, and (b) all three input probes are on. The input power of the probe in each channel is 50 μW. The input power of the control in each channel is 500 μW. The local phase of Ch1 is set to be: $\theta_1 = \phi_1^{(1)} - \phi_2^{(1)} = 0$ and for Ch2: $\theta_2 = \pi$ (if the probe is on). The cell temperature is 60°C.

[S1] W. Cao, X. Lu, X. Meng, J. Sun, H. Shen, and Y. Xiao, Reservoir-mediated quantum correlations in non-Hermitian optical system. Phys. Rev. Lett. 124, 030401 (2020).
[S2] L. Davidovich, Sub-Poissonian processes in quantum optics. Rev. Mod. Phys. 68, 127-173 (1996).
[S3] G. Adesso, and A. Datta, Quantum versus classical correlations in Gaussian states. Phys. Rev. Lett. 105, 030501 (2010).
[S4] P. Giorda, and M. G. A. Paris, Gaussian quantum discord. Phys. Rev. Lett. 105, 020503 (2010).
[S5] I. Novikova, R. L. Walsworth, and Y. Xiao, Electromagnetically induced transparency-based slow and stored light in warm atoms. Laser Photonics Rev. 6, 333 (2012).
FIG. S5. Experimental results with both probes on in Ch1 and Ch2. The transmission probe intensity v.s. the two-photon detuning $\delta_B$, proportional to the applied common magnetic field. Red curve Ch1 is the total transmitted probe power of Ch1. It is the interference results of input probe in Ch1 with transported probe field from Ch2 to Ch1 via the atomic spin transport. Blue curve Ch2 is the transmitted probe power of Ch2 and it is also the result of interference. The local phase of all three channels is set to be: $\theta_0 = 0$, $\theta_1 = 0$ and $\theta_2 = \pi$. The input power of the probe in each channel is 50 $\mu$W. The input power of the control in each channel is 500 $\mu$W. The cell temperature is 60°C. The experiment was performed under the reversed polarization configuration as in Fig. S4.