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Spintronics of non-magnetic chiral media on the example of the Seebeck effect

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Abstract: Spin transport in chiral materials is considered as the main direction for the spintronics development. Experimental work on the study of the spin Seebeck effect shows that in chiral materials it is possible to achieve effective generation of spin-polarized current at distances up to several millimeters. In the presented work, we develop a theoretical approach to describe the spin Seebeck effect based on the spin Hamiltonian for conduction electrons in a chiral medium. Taking into account this Hamiltonian, the equation of electron spin dynamics averaged over the polycrystalline sample is obtained. In the approximation of ideal Fermi-gas and local-quasi-equilibrium distribution, the spin density operator of conduction electrons is constructed. It is shown that averaging over randomly oriented crystallites does not destroy spin ordering at strong spin-orbit interaction, and the temperature gradient generates spin polarization directed predominantly along the temperature gradient. Spin polarization by the described in the paper mechanism does not require external magnetic fields and remanent magnetization and therefore does not interfere with the operation of micro- and nano-sized structures of spintronics. Since the spin Seebeck effect is reciprocal with the spin Peltier effect, the theoretical approach presented in this work can form the basis of new methods for controlling heat flow in spintronics systems.

Keywords: spintronics, spin Seebeck effect, spin-orbit interaction, chiral crystal, polycrystalline conductor

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1. INTRODUCTION

The most important branch of modern spintronics is the generation of spin fluxes in conductors and semiconductors in order to use them in various kinds of microelectronics devices [1]. Conventional methods based on spin-dependent scattering in materials

with strong spin-orbit interaction, such as platinum, provide a ratio of spin-polarized current of less than a percent with coherence lengths on the order of tens of nanometers. This is sufficient for research on spin effects in nanostructures, but not for information and biotechnology. It was shown in [2] that under conditions of inhomogeneous torsion in the steady-state case, the average spin of conduction electrons is oriented predominantly along the charge current density vector. The maximum value of the spin-orbit interaction is observed when the torsion axis is orthogonal to the current density vector and full spin polarization can be achieved. However, this method is energy-consuming. The efficiency of spin polarization drops rapidly with increasing frequency of control actions. This fact also limits the application of controlling spin polarization dynamics by electrical and mechanical actions in modern nanoelectronics.

At the initial stage of spintronics, the spin Seebeck effect (SSE) was considered as a way to efficiently generate spin-polarized current at distances of up to a few millimeters [3]. This effect, consisting in the temperature-gradient-induced generation of spin current, was originally discovered in conducting ferromagnetic metals [4]. The authors explained this effect by magnon and phonon degrees of freedom [5]. Subsequently, the spin Seebeck effect was found in non-magnetic materials [6-8]. The theoretical model of SSE in paramagnetic dielectric material was experimentally confirmed [9]. The microscopic theory of temperature-gradient-driven spin transport in ferromagnets has shown that a temperature gradient is equivalent to an electric field for spin current generation [10]. Similar conclusions are drawn in [11].

It should be noted that the experimentally investigated SSE is essentially a

magneto-thermal-galvanic effect. It is seen in paramagnetic materials in an external magnetic field or in magnetized ferromagnets. The vector of magnetic induction or residual magnetization sets the specific direction in an isotropic matter. In this case, controlling spin polarization by magnetic influence is even slower and more energy consuming than the electrical and mechanical control proposed in [2], where the selected direction sets the torsion vector. A general model of the spin Seebeck effect in non-magnetic materials is currently lacking.

New possibilities for efficient generation of spin currents in spintronics devices are opened by the experimentally detected activation of SSE in nonmagnetic materials by the flow of chiral phonons [8]. According to the authors, chiral phonons possessing angular momentum break the symmetry of the material and create the possibility of spin current generation in the presence of a temperature gradient. Such symmetry breaking exists without external influences in enantiopure chiral crystals such as WSe_2 . Currently, the application of chiral media is considered as the main direction for the development of spintronics [12]. To effectively control large spin current densities the control element must be bulk, rather than thin-film. Therefore, the model of SSE in non-magnetic polycrystalline structures is of interest.

2. MODEL BUILDING

Let us consider the thermally induced polarization of the conduction electron spin in a homogeneous and isotropic polycrystalline metal. The interaction of the collective electron with the crystal field is chosen in the form of a spin-orbit interaction with lattice ions, that is, relativistic corrections in the second order in magnitude $1/c$, where c is the speed of light.

The spin-orbit addition to the electron energy in the field of a crystallite containing N identical ions with effective charge $+Ze$ and coordinates \mathbf{r}_l has the form [13]

$$\hat{V} = \frac{\hbar e}{2m^2 c^2} [\mathbf{E} \times \hat{\mathbf{p}}] \hat{\mathbf{s}}, \quad \mathbf{E}(\mathbf{r}) = -\frac{eZ}{4\pi\epsilon_0} \sum_{l=1}^N \frac{\mathbf{r} - \mathbf{r}_l}{|\mathbf{r} - \mathbf{r}_l|^3}. \quad (1)$$

Here, m is the mass of an electron with charge $-e$, \hbar is the Dirac constant (reduced Planck constant), ϵ_0 is the electric constant. The value of the effective charge Z can be estimated by equating the coordinate of the maximum of the hydrogen-like radial wave function to the covalent radius of the atom. For example, for platinum the atomic radius is $1.39 \cdot 10^{-10} m$, which for a 6s shell corresponds to $Z \approx 22.45$.

Let us construct the spin Hamiltonian of the perturbation (1) for the conduction electron by averaging it over the coordinates [14]. The wave function of the conduction electron with wave vector \mathbf{k} is expressed as a Wannier function [15]

$$\psi_{\mathbf{k}}(\mathbf{r}) = \frac{1}{\sqrt{N}} \sum_{n=1}^N \Psi(\mathbf{r} - \mathbf{R}_n) \exp(i\mathbf{k}\mathbf{R}_n),$$

Here, $\Psi(\mathbf{r})$ is the atomic function of the electron, \mathbf{R}_n is the translation vector. Replacing the variables $\mathbf{r} - \mathbf{r}_l \rightarrow \mathbf{r}$ we obtain

$$\hat{V}_{\mathbf{k}} = -\frac{\hbar^2 e^2 Z \hat{\mathbf{s}}_{\mathbf{k}}}{8\pi\epsilon_0 m^2 c^2 N} \sum_{n,m,l=1}^N \exp(i\mathbf{k}(\mathbf{R}_n - \mathbf{R}_m)) \times \langle \Psi(\mathbf{r} + \mathbf{r}_l - \mathbf{R}_m) | \frac{\hat{\mathbf{1}}}{r^3} | \Psi(\mathbf{r} + \mathbf{r}_l - \mathbf{R}_n) \rangle.$$

In the nearest neighbor approximation we retain in the right-hand side only the summands for which $\mathbf{R}_n - \mathbf{r}_l = \mathbf{R}_m - \mathbf{r}_l = 0$, or $\mathbf{R}_n - \mathbf{r}_l = \mathbf{a}_\nu$ and $\mathbf{R}_m - \mathbf{r}_l = 0$, or $\mathbf{R}_n - \mathbf{r}_l = 0$ and $\mathbf{R}_m - \mathbf{r}_l = \mathbf{a}_\nu$, where \mathbf{a}_ν is the vector drawn to the nearest neighbor node. In the first order of smallness in $\mathbf{k}\mathbf{a}_\nu$ we obtain

$$p_{ij} = \begin{bmatrix} \cos(\alpha)\cos(\gamma) - \sin(\alpha)\cos(\beta)\sin(\gamma) & -\cos(\alpha)\sin(\gamma) - \sin(\alpha)\cos(\beta)\cos(\gamma) & \sin(\alpha)\sin(\beta) \\ \sin(\alpha)\cos(\gamma) + \cos(\alpha)\cos(\beta)\sin(\gamma) & -\sin(\alpha)\sin(\gamma) + \cos(\alpha)\cos(\beta)\cos(\gamma) & -\cos(\alpha)\sin(\beta) \\ \sin(\beta)\sin(\gamma) & \sin(\beta)\cos(\gamma) & \cos(\beta) \end{bmatrix},$$

$$\hat{V}_{\mathbf{k}} = \hat{\mathbf{s}}_{\mathbf{k}} (\mathbf{I}_0 + \mathbf{I}_1 + \mathbf{I}_{\mathbf{k}}), \quad \mathbf{I}_0 = -\frac{\hbar^2 e^2 Z}{8\pi\epsilon_0 m^2 c^2} \langle \Psi | \frac{\hat{\mathbf{1}}}{r^3} | \Psi \rangle,$$

$$\Psi_\nu(\mathbf{r}) = \Psi(\mathbf{r} + \mathbf{a}_\nu), \quad \mathbf{I}_1 = -\frac{\hbar^2 e^2 Z}{4\pi\epsilon_0 m^2 c^2} \text{Re} \sum_\nu \langle \Psi_\nu | \frac{\hat{\mathbf{1}}}{r^3} | \Psi \rangle, \quad (2)$$

$$\mathbf{I}_{\mathbf{k}} = -\frac{\hbar^2 e^2 Z}{4\pi\epsilon_0 m^2 c^2} \text{Im} \sum_\nu (\mathbf{k}\mathbf{a}_\nu) \langle \Psi_\nu | \frac{\hat{\mathbf{1}}}{r^3} | \Psi \rangle.$$

For s -electrons the vectors \mathbf{I}_0 , \mathbf{I}_1 , $\mathbf{I}_{\mathbf{k}}$ are zero. The conduction bands overlap in metals and part of the conduction electrons can be formed by collectivization of p -electrons.

Let us consider the macroscopic area of a polycrystalline metal. For any state of the electron one can choose the direction of the quantization axis ($z_{\tilde{\chi}}$ -axis) so that the projection of its orbital momentum on this axis has a certain value $l_{\tilde{\chi}} = l$. The electron energy in an atom under an electric field depends on the projection of its orbital momentum on the direction of the field [14]. Therefore, the orientation of atomic orbitals is determined by the orientation of the crystallo-physical axes of the crystallite, and we can consider that the vectors \mathbf{I}_0 and \mathbf{I}_1 , in relations (2) are written in the coordinate system associated with the symmetry axes of the crystallite.

Let us introduce the laboratory coordinate system associated with the instruments that set the conduction current and measure the spin components. Therefore, the wave vector and spin vector of the conduction electrons should be considered as specified in the laboratory coordinate system. The components of vectors and tensors in the laboratory system will be denoted by non-primed indices, but in the coordinate system, associated with the crystal axes, by primed these.

The vectors \mathbf{I}_0 and \mathbf{I}_1 are converted to the laboratory coordinate system as follows $I_{\alpha'} = p_{\alpha\alpha'} \cdot I_{\alpha}$, where $p_{\alpha\alpha'}$ is the unitary rotation matrix. It is convenient to express the rotation matrix through Euler angles:

Here, $0 \leq \alpha \leq 2\pi$ is the precession angle, $0 \leq \beta \leq \pi$ is the nutation angle, $0 \leq \gamma \leq 2\pi$ is the angle of proper rotation. Then for a macroscopically isotropic conductor averaging over random crystallite orientations leads to averaging over random uniformly distributed Euler angles.

Let us average the vectors \mathbf{I}_0 and \mathbf{I}_1 over random crystallite orientations. In an isotropic polycrystal $\langle \mathbf{I}_0 \rangle = \langle \mathbf{I}_1 \rangle = 0$. For a symmetric crystal in which each nearest neighbor with vector \mathbf{a}_ν corresponds to a neighbor with vector $\mathbf{a}_{-\nu} = -\mathbf{a}_\nu$, the last relation of formula (2) takes the form

$$\mathbf{I}_k = -\frac{\hbar^2 e^2 Z}{4\pi\epsilon_0 m^2 c^2} \text{Im} \sum_{\nu>0} (\mathbf{k}\mathbf{a}_\nu) \langle \Psi_\nu - \Psi_{-\nu} | \frac{\hat{\mathbf{I}}}{r^3} | \Psi \rangle. \quad (3)$$

In a symmetric crystal, the function $\Psi_\nu - \Psi_{-\nu}$ has parity opposite to the parity of the function Ψ . Therefore, all summands in (6) for the undeformed crystal are zero.

At distortion $r'_\alpha = r_\alpha + u_\alpha(\mathbf{r})$ we obtain in (3) $\Psi(\mathbf{r}') = \Psi(\mathbf{r}) + r_\beta \frac{\partial \Psi}{\partial r_\alpha} \frac{\partial u_\alpha}{\partial r_\beta}$, $\hat{I}'_\alpha = \hat{I}_\alpha - i\epsilon_{\alpha\beta\gamma} \left(u_\beta \frac{\partial}{\partial r_\gamma} - r_\beta \frac{\partial u_\delta}{\partial r} \frac{\partial}{\partial r_\delta} \right)$. Here, $\epsilon_{\alpha\beta\gamma}$ is the Levi-Civita tensor. For inhomogeneous torsional strain along the \mathbf{n} axis of the form $\boldsymbol{\Omega}(\mathbf{r}) = \mathbf{n}(\mathbf{r}\mathbf{n})\omega$, we obtain

$$I_{k\alpha} = -\frac{\hbar^2 e^2 Z \omega}{4\pi\epsilon_0 m^2 c^2} \epsilon_{\alpha\beta\gamma} n_\beta n_\delta k_\sigma a_{\nu\sigma} \times \text{Im} \langle \Psi_\nu - \Psi_{-\nu} | \frac{r_\gamma \hat{I}'_\delta}{r^3} | \Psi \rangle, \quad (4)$$

Here, the summation over ν by nearest neighbor pairs is implied.

3. SPINTRONICS OF CHIRAL MEDIA

The product of the value ω by the distance \mathbf{a}_ν to the nearest neighbor in the direction of the torsion axis \mathbf{n} characterizes the torsion of the unit cell, i.e., the rotation angle of its crystallographic planes relative to its neighbors. Such torsion exists in chiral

crystals such as WSe₂ and layered hybrid perovskites [10]. Therefore, it can be considered that vector \mathbf{n} , as well as vectors \mathbf{a}_ν are set in the system of crystallophysical axes, and vector \mathbf{k} – in the laboratory system. Then formula (4) can be written in the form

$$\langle I_{k\alpha} \rangle = -\frac{\hbar^2 e^2 Z \omega}{4\pi\epsilon_0 m^2 c^2} \langle p_{\alpha\alpha} p_{\sigma\sigma}^{-1} \rangle \epsilon_{\alpha\beta\gamma} n_\beta n_\delta k_\sigma a_{\nu\sigma} \times \text{Im} \langle \Psi_\nu - \Psi_{-\nu} | \frac{r_\gamma \hat{I}'_\delta}{r^3} | \Psi \rangle. \quad (5)$$

Let us denote $\mathbf{s}_k(t) = \langle \hat{\mathbf{s}}_k(t) \rangle = \text{Sp}(\hat{\mathbf{s}}_k(t) \hat{\rho}(t))$, where $\hat{\rho}(t)$ is the density operator. Then

$$\frac{d\mathbf{s}_k}{dt} = \left\langle \frac{d\hat{\mathbf{s}}_k}{dt} \right\rangle + \text{Sp} \left(\hat{\mathbf{s}}_k \frac{d\hat{\rho}}{dt} \right). \quad (6)$$

The dynamics of the conduction electron spin caused by the perturbation (2) is described by equation [14]

$$\begin{aligned} \frac{d\hat{\mathbf{s}}_{k\alpha}}{dt} &= \frac{i}{\hbar} [\hat{V}_k, \hat{s}_\alpha] = \epsilon_{\alpha\beta\gamma} I_{k\beta} \hat{s}_{k\gamma}, \\ \left\langle \frac{d\hat{\mathbf{s}}_k}{dt} \right\rangle &= [\langle \mathbf{I}_k \rangle \times \mathbf{s}_k] + \langle [\delta \mathbf{I}_k \times \delta \hat{\mathbf{s}}_k] \rangle, \\ \delta \mathbf{I}_k &= \mathbf{I}_k - \langle \mathbf{I}_k \rangle, \quad \delta \hat{\mathbf{s}}_k = \hat{\mathbf{s}}_k - \langle \hat{\mathbf{s}}_k \rangle. \end{aligned} \quad (7)$$

In the steady-state, when $d\hat{\rho}/dt = 0$, only the first summand in the righthand side of equation (6) remains and equation (7) takes the form

$$\frac{d\mathbf{s}_k}{dt} = [\langle \mathbf{I}_k \rangle \times \mathbf{s}_k] + \langle [\delta \mathbf{I}_k \times \delta \hat{\mathbf{s}}_k] \rangle. \quad (8)$$

The modulus of the average spin is not conserved because of the second summand in the right-hand side of (8). Therefore, it can be considered as relaxational and written in the form $-(\mathbf{s}_k - \mathbf{s}_k^e)/\tau$. Here, \mathbf{s}_k^e is the equilibrium value of the sample average spin, τ is the longitudinal relaxation time.

In the absence of influences, a quasi-equilibrium distribution is established with a density operator [16]

$$\hat{\rho}^q(t) = \exp \left\{ -\Phi(t) - \int_V \theta(t, \mathbf{r}) \hat{h}_k(t, \mathbf{r}) d^3r \right\}, \quad (9)$$

$$\Phi(t) = \ln \text{Sp} \exp \left\{ - \int_V \theta(t, \mathbf{r}) \hat{h}_k(t, \mathbf{r}) d^3r \right\}.$$

Here, $\theta(t, \mathbf{r}) = 1/(k_B T(t, \mathbf{r}))$, k_B is the Boltzmann constant, $T(t, \mathbf{r})$ is the local temperature, $\hat{h}_k(t, \mathbf{r})$ is the density operator of the Hamiltonian, satisfying the relations [17]

$$\hat{H}_k(t) = \int_V \hat{h}_k(t, \mathbf{r}) d^3r, \quad \frac{\partial \hat{h}_k(t, \mathbf{r})}{\partial t} = -\text{div} \hat{\mathbf{q}}_k(t, \mathbf{r}), \quad (10)$$

Here, $\hat{\mathbf{q}}_k(t, \mathbf{r})$ is the flux density operator of the Hamiltonian.

We used the ideal Fermi-gas approximation for conduction electrons. The applicability of this model to conduction electrons in metals is justified by the fact that the thermodynamics of the Fermi-system is determined by its microscopic structure only near the Fermi-surface and is completely independent of what happens outside the range of order $k_B T$. As a result, the denser the Fermi-gas in the metal, the more ideal it is [18]. Experimental studies of the temperature dependence of the electron heat capacity in metals show that it corresponds well to the model of an ideal Fermi-gas with a scalar effective mass m^* . For most metals, the effective mass of the conduction electron is close to the mass of the free electron $m^* \approx m$. Such a model allows us to describe the collective of electrons with wave vector \mathbf{k} by its density operator and put $\hat{\mathbf{q}}_k(t, \mathbf{r}) = \hbar \mathbf{k} \hat{h}_k(t, \mathbf{r})/m^*$. Then, from (9) and (10) we obtain

$$\frac{d\hat{\rho}^q}{dt} = -\hat{\rho}^q \left\{ \frac{d\Phi}{dt} + \int_V \frac{\partial \theta}{\partial t} \hat{h}_k d^3r - \int_V \text{div}(\theta \hat{\mathbf{q}}_k) d^3r + \int_V \hat{\mathbf{q}}_k \text{grad} \theta d^3r \right\}. \quad (11)$$

Analytical averaging of equation (5) gives $\langle \mathbf{I}_k \rangle = \mathcal{J} \mathbf{k}$, where

$$J = \frac{\hbar^2 e^2 Z \omega}{24 \pi \epsilon_0 m^2 c^2} \text{Im} \langle \Psi_\nu - \Psi_{-\nu} | \frac{(\mathbf{n}\hat{\mathbf{l}})[\mathbf{a}_\nu \times \mathbf{n}]\mathbf{r}}{r^3} | \Psi \rangle.$$

Let us substitute the ratio (11) into the second summand in the right-hand side of equation (6), restricting to summands proportional to $\partial \theta / \partial \mathbf{r}$. Taking into account (10), we obtain

$$\frac{d\mathbf{s}_k}{dt} = J [\mathbf{k} \times \mathbf{s}_k] - \frac{\mathbf{s}_k - \mathbf{s}_k^e}{\tau} - \frac{\hbar}{m^*} \langle \hat{\mathbf{s}}_k \hat{H}_k \rangle^e \mathbf{k} \frac{\partial \theta}{\partial \mathbf{r}} \quad (12)$$

In the absence of a magnetic field, the unperturbed Hamiltonian of an ideal Fermi-gas does not depend on spin variables, so

$$\langle \hat{\mathbf{s}}_k \hat{H}_k \rangle^e = \langle \hat{\mathbf{s}}_k \rangle^e \langle \hat{H}_k \rangle^e = \mathbf{s}_k^e H(\mathbf{k}).$$

Let us introduce the spin density operator of conduction electrons

$$\hat{\mathbf{s}}(t) = \int_V \hat{\mathbf{s}}(t, \mathbf{r}) d^3r,$$

$$\hat{\mathbf{s}}(t, \mathbf{r}) = \frac{1}{8\pi^3} \int \hat{\mathbf{s}}_k(t) f(\mathbf{k}, \mathbf{r}) d^3k, \quad (13)$$

$$f(\mathbf{k}, \mathbf{r}) = \left[\exp((H(\mathbf{k}) - F(\mathbf{r}))\theta(\mathbf{r})) + 1 \right]^{-1}.$$

For an isotropic Fermi-gas $H(\mathbf{k}) = \hbar^2 k^2 / (2m^*)$. The Fermi energy in (13) is found as $F(\mathbf{r}) = \hbar^2 k_F^2(\mathbf{r}) / (2m^*)$, $k_F^3(\mathbf{r}) = 6\pi^2 n_e(\mathbf{r})$. Here, $n_e(\mathbf{r})$ is the conduction electron density. The first summand in the right-hand side of equation (12) is zero when averaging of the form (13).

The steady-state in (8) corresponds to the orientation of \mathbf{s}_k along the vector \mathbf{k} . Averaging the perturbation (1) by quantum state and by random orientations, we obtain that the addition to the energy of states when spin is oriented parallel or antiparallel to the vector \mathbf{k} is $\pm \hbar J k / 2$. Then

$$\mathbf{s}^e = \frac{1}{16\pi^3} \int \frac{\mathbf{k}}{k} \left\{ f \left(H(k^2) - \frac{\hbar J k}{2} \right) - f \left(H(k^2) + \frac{\hbar J k}{2} \right) \right\} d^3k = 0.$$

The wave vector average of the last summand in the right part (12) is equal to

$$\begin{aligned} & \frac{\hbar^2 J}{16\pi^3 m^*} \int k \frac{df}{dH} \frac{\mathbf{k}}{k} H(k^2) \left(\mathbf{k} \frac{\partial \theta}{\partial \mathbf{r}} \right) d^3 k = \\ & = \frac{3F(\mathbf{r})J}{2\pi} n_e(\mathbf{r}) \frac{\partial \theta}{\partial \mathbf{r}}, \end{aligned}$$

and the steady-state spin polarization is parallel to the temperature gradient

$$\mathbf{s}(\mathbf{r}) = \frac{3F(\mathbf{r})J\tau}{2\pi k_B} \frac{n_e(\mathbf{r})}{T^2} \frac{\partial T}{\partial \mathbf{r}}. \quad (14)$$

4. CONCLUSION

Formula (14) describes the well-known longitudinal spin Seebeck effect, when the direction of spin polarization is determined by the temperature gradient even in the absence of a charge current. Of course, the temperature gradient will also create the classical (charge) Seebeck effect in the conductor. Therefore, if a conductor is connected in a closed circuit, a charge current with longitudinal spin polarization will appear in it. The charge current in the conductor can also be generated by an external current source, which opens up the possibility of rapid electrical control of spin polarization. The high efficiency of such control is due to the fact that in chiral media the torsion of the unit cell can reach several radians, which is many orders of magnitude more than can be obtained by mechanical stresses [2] without irreversible destruction. The possibility of spin polarization generation at a distance of several millimetres in polycrystalline chiral media was shown in [19].

Spin polarization by the mechanism described in this work does not require external magnetic fields and residual magnetization and therefore does not interfere with the work in micro- and nano-sized spintronics structures. The energy required for spin

polarization by the proposed mechanism comes from the heat fluxes in the medium due to the temperature gradient. However, the functioning of spintronics devices, as well as any microelectronics devices, is always accompanied by the generation of heat fluxes and temperature gradients. In this sense, the energy expended for spin polarization by the proposed mechanism in spintronics systems is "free".

In addition, the Seebeck spin effect is reciprocal with the Peltier spin effect. Reciprocity relations for them, at least in magnetic dielectrics, have been experimentally verified [20]. General reciprocity relations between Peltier and Seebeck spin effects, including for conducting media in the presence of mechanical strain, were obtained in [21]. This suggests that the dynamical effect (12) presented in this work may form the basis for new methods of heat flow control in spintronics systems.

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