



Perspective Article

Spin caloritronics, origin and outlook

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ABSTRACT

Spin caloritronics refers to research efforts in spintronics when a heat current plays a role. In this review, we start out by reviewing the predictions that can be drawn from the thermodynamics of irreversible processes. This serves as a conceptual framework in which to analyze the interplay of charge, spin and heat transport. This formalism predicts tensorial relations between vectorial quantities such as currents and gradients of chemical potentials or of temperature. Transverse effects such as the Nernst or Hall effects are predicted on the basis that these tensors can include an anti-symmetric contribution, which can be written with a vectorial cross-product. The local symmetry of the system may determine the direction of the vector defining such transverse effects, such as the surface of an isotropic medium. By including magnetization as state field in the thermodynamic description, spin currents appear naturally from the continuity equation for the magnetization, and dissipative spin torques are derived, which are charge-driven or heat-driven. Thermodynamics does not give the strength of these effects, but may provide relationships between them. Based on this framework, the review proceeds by showing how these effects have been observed in various systems. Spintronics has become a vast field of research, and the experiments highlighted in this review pertain only to heat effects on transport and magnetization dynamics, such as magneto-thermoelectric power, or the spin-dependence of the Seebeck effect, the spin-dependence of the Peltier effect, the spin Seebeck effect, the magnetic Seebeck effect, or the Nernst effect. The review concludes by pointing out predicted effects that are yet to be verified experimentally, and in what novel materials the standard thermal spin effects could be investigated.

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1. Introduction

The term “spin caloritronics” was coined to refer to all transport phenomena that involve spin and heat [1,2]. This sub-field of spintronics has drawn considerable attention since 2009, when an international workshop launched this emerging field. Specialists have met annually ever since, as this theme of research attracts more and more researcher. 2016 has seen the seventh “Spin Caloritronics” workshop [3]. The German Physical Society (DFG) launched a priority program to foster research in this field [4]. Recently, the US Department of Energy funded a UC-Riverside-based research cluster “SHINES” (Spins and Heat in Nanoscale Electronic Systems) consisting 14 research groups from 7 institutions from coast to coast. Thus, spin caloritronics as an interdisciplinary field of magnetism, thermoelectrics and microelectronics will continuously attract global interests of cutting-edge scientific research [5].

This review is structured as follows. The definition of some of the basic concepts of spin caloritronics are presented within a thermodynamic framework. Then, experiments on spin-dependent transport phenomena are reported, provided that the effects under study were driven by a heat current. Others have also attempted to review this field, which is bursting in all kinds of new directions [6]. This paper does not seek to analyze the theoretical efforts carried out to account quantitatively for spin caloritronics phenomena.

A few examples can be mentioned to highlight the diversity of approaches. The magneto-thermoelectric power (MTEP) of magnetic tunnel junctions (MTJ) has been modeled using the Buttiker–Landauer formalism [7]. The key material properties of the magnetic electrodes has been identified by ab initio calculations [8]. The details of spin transport across tunnel barriers have also been calculated [9,10]. Dissipative torques may affect magnetization dynamics, in conductors as well as in metals. These are often known as non-adiabatic torques or entropic torques. This type of torque, associated with the Dzyaloshinskii–Moriya interaction, was recently identified and its role in the dynamics of

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domain walls was analyzed [11]. The method of master equations was used to describe a single quantum dot with spin-dependent electron temperatures due to its connection to ferromagnetic leads [12]. The Kubo formalism of linear response theory has been applied to describe thermal spin orbit torques and the reciprocal effect, the heat current associated with magnetization dynamics [13,14]. The existence of thermal spin orbit torques of electronic origin has been predicted using a Berry phase description [15]. magnonic contributions to the spin orbit torque have also been identified [16,17] including thermally-driven torques [18]. The effect of thermally-driven torque on domain wall motion has been simulated by using atomistic spin dynamics [19], or by considering the drift caused by thermal fluctuations in a temperature gradient [20]. The phenomenon that has boosted the interest for spin caloritronics is the spin Seebeck effect. It has been accounted for by considering the dynamics of magnons in a temperature gradient [21,22]. The Boltzmann transport theory was also used to describe this effect [23,24]. A key consideration to understand this and other related effects is the non-equilibrium condition that may arise near an interface [25]. In particular, magnons and phonons may have different temperatures [26]. The accumulation of magnons at an interface, due to presence of a temperature gradient, has been described in terms of a Bose-Einstein condensation [27,28]. A review on the theory of spin Seebeck effect can be found in [29].

1.1. Thermodynamic description of transport phenomena, definitions

In this section, we first recall that the thermodynamics of irreversible processes provides a framework in which to define Ohm's law [30] and the Hall effect [31], and in a similar fashion, the Seebeck and the Nernst effect. The transverse effects (Hall and Nernst) are expected merely on the basis of a symmetry argument. One should keep in mind that thermodynamics is a powerful way of establishing relations among various physical quantities, but thermodynamics does not provide quantitative estimates for any of the transport coefficients, as it does not address the underlying micro-mechanisms responsible for the transport phenomena that it predicts.

In thermodynamics, diffusive current densities are defined for each density of extensive variables. For example, entropy is an extensive variable and the current density \mathbf{j}_s is defined. A lengthy development, which takes into account Newton's law and the first principle of thermodynamics, shows that the density of entropy source ρ_s is given by an expression of the generic form [32–34]:

$$\rho_s = \frac{1}{T} \sum_i \mathbf{j}_i \cdot \mathbf{F}_i \quad (1)$$

The generalized forces \mathbf{F}_i are the gradients of the intensive variables conjugated to the extensive state variables numbered by the index $i = 1..n$. Here we consider $(s, \{n_A\})$, the state variable densities for entropy and quantities of substance A, B, \dots , respectively. For a substance A , the associated generalized force is the gradient of its electrochemical potential, $\bar{\mu}_A = \mu_A + q_A V$, where V is the electrostatic potential and q_A the elementary charge of substance A . Thus, in the following, $\mathbf{F}_A = \nabla \mu_A + q_A \nabla V$. The generalized force associated with entropy is $-\nabla T$.

The second principle of thermodynamics imposes that $\rho_s \geq 0$. This implies the following constitutive equations:

$$\begin{cases} \mathbf{j}_s = L_{ss} \cdot (-\nabla T) + \sum_B L_{sB} \cdot \mathbf{F}_B \\ \mathbf{j}_A = L_{As} \cdot (-\nabla T) + \sum_B L_{AB} \cdot \mathbf{F}_B \end{cases} \quad (2)$$

The Onsager matrix that underlies (2) must be positive definite in order to satisfy the condition $\rho_s \geq 0$. In all generality, the Onsager matrix elements such as L_{ss} , L_{sA} and L_{AA} are tensors. The

Onsager–Casimir relations [35,36] imposes the following symmetry on the coefficients of (2):

$$L_{\alpha\beta}(s, \{n_A\}, \mathbf{B}) = \varepsilon_\alpha \varepsilon_\beta L_{\beta\alpha}(s, \{n_A\}, -\mathbf{B}) \quad (3)$$

where, according to a result derived from statistical physics, the parameters $\varepsilon_\alpha = \pm 1$, $\varepsilon_\beta = \pm 1$. The parameters ε_α and ε_β are positive if the corresponding generalized forces \mathbf{F}_α and \mathbf{F}_β are invariant under time reversal, and they are negative in the opposite case. For the vectorial quantities considered here and in the following sections, the parameters ε_α and ε_β are positive because the corresponding generalized vectorial forces \mathbf{F}_α and \mathbf{F}_β are invariant under time reversal.

We consider now the particular case of an isotropic medium composed of one substance A , in the absence of chemical effect ($\mu_A = 0$), at a homogeneous temperature T but with an electric field applied to the system. The result (2) implies the tensorial expression of Ohm's law is:

$$\mathbf{j}_A = \frac{-1}{q_A} \sigma \cdot \nabla V \quad (4)$$

In all generality, the tensor σ can be decomposed in a symmetric tensor and an antisymmetric tensor, $\sigma = \sigma^s + \sigma^a$. For the symmetric part, we will simply consider the isotropic case $\sigma^s = \sigma \mathbb{1}$. The antisymmetric part provides richer physical insight. We can write σ^a in the form of a vectorial product, $\sigma^a \cdot \mathbf{x} = \sigma_\perp \hat{\mathbf{u}} \times \mathbf{x}$, for any vector \mathbf{x} [37]. The particular physical properties of the system determine the direction and modulus of the vector $\sigma_\perp \hat{\mathbf{u}}$. Thus, the decomposition of the conductivity tensor into symmetric and antisymmetric parts allows us to write:

$$\mathbf{j}_A = -\frac{1}{q_A} \sigma \nabla V - \frac{\sigma_\perp}{q_A} (\hat{\mathbf{u}} \times \nabla V) \quad (5)$$

In particular, if a magnetic induction field \mathbf{B} is applied to the system, the Onsager–Casimir relations (3) imply that $\sigma_{ij}^a(\mathbf{B}) = \sigma_{ji}^a(-\mathbf{B}) = -\sigma_{ji}^a(\mathbf{B})$, where the last equality is the expression of the antisymmetry of the tensor. Therefore, σ^a is an antisymmetric function of \mathbf{B} . For an isotropic medium, this means that $\hat{\mathbf{u}}$ must be in the direction of \mathbf{B} . Hence, (4) includes the Hall effect:

$$\mathbf{j}_A = -\frac{1}{q_A} \sigma \nabla V - \frac{\sigma_\perp}{q_A} (\hat{\mathbf{B}} \times \nabla V) \quad (6)$$

Furthermore, if we are near the surface of the medium, the isotropy breaks down and we have an axial symmetry about the normal $\hat{\mathbf{n}}$ to the surface. It is easy to show that the only way to have (5) with the symmetry about $\hat{\mathbf{n}}$ is to write:

$$\mathbf{j}_A = -\frac{1}{q_A} \sigma \nabla V - \frac{\sigma_\perp}{q_A} (\hat{\mathbf{n}} \times \nabla V) \quad (7)$$

Now, we consider transport effects due to the presence of a temperature gradient and in the absence of an electric field ($\nabla V = \mathbf{0}$). From (2), we deduce:

$$\mathbf{j}_A = \frac{-1}{q_A} \sigma \cdot \boldsymbol{\varepsilon} \cdot \nabla T \quad (8)$$

The notation has been chosen to keep track of the Seebeck effect and Ohm's law. Hence, we define $\sigma = q_A^2 L_{AA}$ and $\boldsymbol{\varepsilon} = (1/q_A) L_{AA}^{-1} \cdot L_{As}$. The tensor $\boldsymbol{\varepsilon}$ can be decomposed in symmetric and antisymmetric parts. Following what we did for the conductivity tensor, we write $\boldsymbol{\varepsilon} \cdot \mathbf{x} = \varepsilon \mathbf{x} + \varepsilon_\perp (\hat{\mathbf{u}} \times \mathbf{x})$. Then (8) becomes:

$$\mathbf{j}_A = -\frac{\sigma \varepsilon}{q_A} \nabla T - \frac{\sigma \varepsilon_\perp + \sigma_\perp \varepsilon}{q_A} (\hat{\mathbf{u}} \times \nabla T) \quad (9)$$

Just as discussed regarding the conductivity, the unit vector $\hat{\mathbf{u}}$ can designate either the direction of the magnetic induction field or the normal to the surface.

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