Room-temperature spin-orbit torque in NiMnSb

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Materials that crystallize in diamond-related lattices, with Si and GaAs as their prime examples, are at the foundation of modern electronics. Simultaneously, inversion asymmetries in their crystal structure and relativistic spin-orbit coupling led to discoveries of non-equilibrium spin-polarization phenomena that are now extensively explored as an electrical means for manipulating magnetic moments in a variety of spintronic structures. Current research of these relativistic spin-orbit torques focuses primarily on magnetic transition-metal multilayers. The low-temperature diluted magnetic semiconductor (Ga,Mn)As, in which spin-orbit torques were initially discovered, has so far remained the only example showing the phenomenon among bulk non-centrosymmetric ferromagnets. Here we present a general framework, based on the complete set of crystallographic point groups, for identifying the potential presence and symmetry of spin-orbit torques in non-centrosymmetric crystals. Among the candidate room-temperature ferromagnets we chose to use NiMnSb, which is a member of the broad family of magnetic Heusler compounds. By performing all-electrical ferromagnetic resonance measurements in single-crystal epilayers of NiMnSb we detect room-temperature spin-orbit torques generated by effective fields of the expected symmetry and of a magnitude consistent with our ab initio calculations.

Recently, the focus of spintronic research has been on phenomena that exploit the relativistic transfer between the linear momentum of the electron and spin. The spin Hall effect (SHE) and the inverse spin galvanic effect (ISGE), experimentally discovered more than a decade ago as companion phenomena in GaAs-based structures, play a fundamental role in this so-called spin–orbitronics research field. Unlike the SHE, generating a uniform spin current and a resulting surface/interface spin accumulation, the ISGE induces directly a uniform non-equilibrium spin polarization in spin-orbit-coupled crystals with broken inversion symmetry. The ISGE (also called the Edelstein effect) was experimentally discovered in GaAs, whose bulk zinc-blende unit cell is non-centrosymmetric.

The non-equilibrium spin polarization generated by the ISGE and the corresponding effective field can induce spin-orbit torques on the magnetization in non-centrosymmetric magnetic crystals, as demonstrated in the low-Curie-temperature diluted magnetic semiconductor (Ga,Mn)As (refs 8–11), or in magnetic multilayers with structural inversion asymmetry (ref. 12). The multilayers, typically comprising an interface of a high-Curie-temperature transition-metal ferromagnet and a strongly spin-orbit-coupled paramagnet, have attracted most of the attention so far for their direct relevance to magnetic memory and other spintronic applications. However, the single-layer magnets we focus on in this work have their own merits in the research of current-induced spin–orbit torques. The effects in these systems originate from the ISGE and corresponding broken symmetries of bulk crystals, whereas in the magnetic multilayers the ISGE mechanism is entangled with the SHE (refs 13–16). The understanding of spin–orbit torques in bulk non-centrosymmetric crystals is therefore more straightforward, and their potentially utility in spintronic devices is more robust against unintentional disorder and structural imperfections than the spin–orbit coupling phenomena generated within a few atomic planes near the ferromagnet/paramagnet interface.

As the low-Curie-temperature disordered alloy (Ga,Mn)As has so far been the only material studied in which spin–orbit torques are generated by the bulk ISGE, we start this paper with a general analysis of the potential presence and of the symmetries of the ISGE-induced spin–orbit fields considering all 21 crystal point groups with broken inversion symmetry. This opens the route for exploring and potentially exploiting spin–orbit torques in non-centrosymmetric crystals beyond the singular and academic example of (Ga,Mn)As. We then identify NiMnSb as a candidate ferromagnet for observing the spin–orbit torque in an ordered magnetic crystal at room temperature.

NiMnSb is a member of the broad family of magnetic Heusler compounds which in bulk is predicted to be a half-metal ferromagnet and has a Curie temperature of 730 K (refs 17, 18). The material can have a low Gilbert damping constant $\sim 10^{-3}$ and tunable magnetic anisotropies when grown in thin films (ref. 19). For these characteristics, NiMnSb has been used in earlier spintronics studies of non-relativistic magneto-transport and spin dynamics effects based on spin angular momentum transfer between carriers and magnetization in magnetic-multilayer devices (ref. 20). In the second part of our paper we report experiments in single-crystal epilayers of NiMnSb in which we identify the ISGE-induced spin–orbit torques by employing the all-electrical ferromagnetic resonance (FMR) technique (ref. 21). Experimental results are compared with ab initio transport theory calculations of the spin–orbit torque in the studied NiMnSb material.

Symmetry of spin–orbit fields

In Fig. 1a,b we illustrate examples of the relativistic non-equilibrium spin polarizations that occur in the family of diamond-related...
lattices for the case when the two inversion-partner lattice sites of the unit cell are occupied by the same or by different atomic species. The ISGE responsible for these polarization effects requires the spin–orbit coupling to be combined with inversion asymmetries in the crystal structure\textsuperscript{24,25}. Each of the two atomic lattice sites in the unit cell of the crystals shown in Fig. 1a,b have an inversion-asymmetric local environment\textsuperscript{26} which allows the generation of the local current-induced spin polarization\textsuperscript{27,28}.

By symmetry, the global polarization vanishes when integrated over the entire unit cell if the two sites are occupied by the same atom, because the unit cell has an inversion centre (highlighted by a red dot in Fig. 1a). Only local current-induced polarizations of the same magnitude and opposite sign at the two inversion-partner sites remain non-zero in this case (red arrows in Fig. 1a). When the sites are occupied by different atoms, the unit cell is globally inversion asymmetric allowing the generation of a net global spin polarization with a non-zero integral value over the unit cell (red arrow in Fig. 1b). For completeness we point out that in these diamond-related cubic lattices an additional symmetry lowering has to be introduced to allow for the above non-equilibrium polarization effects. For example, in thin films with tetragonal distortion due to a substrate lattice-matching strain, the current-induced spin polarization acquires a Dresselhaus symmetry with respect to the crystal direction of the applied electric field\textsuperscript{9,10}, as illustrated in Fig. 1c. A shear strain, on the other hand, generates spin polarization with the Rashba symmetry\textsuperscript{10,11} (see Fig. 1c).

In Table 1 we summarize symmetries of the current-induced spin polarizations and the corresponding spin–orbit fields, $\mathbf{h} = \chi \mathbf{E}$, based on the analysis of all crystals belonging to the 21 point groups with broken inversion symmetry. Here $\mathbf{E}$ is the applied electric field and $\chi$ is the response tensor. To find this tensor for each point group we made use of Neumann’s principle—that is, we looked for axial tensors that satisfy $\chi = \det(R)R\chi R^{-1}$ for all symmetry operations $R$ (see Supplementary Section 4). The resulting current-induced spin–orbit field symmetries with respect to the applied electric field direction are of the generalized Rashba or Dresselhaus form shown in Fig. 1c, or the spin–orbit field is parallel to $\mathbf{E}$, or the symmetries are a combination of these three forms. Note that we consider here only the spin–orbit field terms that are independent of the magnetization and that dominate in the NiMnSb samples studied below. For the discussion of higher-order terms that depend on magnetization, see the Methods.

The cubic half-Heusler lattice of NiMnSb, shown in Fig. 1d, shares the $-43m$ point-group symmetry with the cubic diamond
Table 1 | Current-induced spin–orbit field symmetries for all point groups with broken inversion symmetry.

| Point group | Field-like $\chi$ | Point group | Field-like $\chi$ |
|-------------|-------------------|-------------|-------------------|
| 2           | $(x_{11} 0 0)$    | 312         | $(x_{11} 0 0)$    |
|             | $(0 x_{22} 0)$    |             | $(0 x_{11} 0)$    |
|             | $(x_{31} 0 x_{33})$ |             | $(0 0 x_{33})$    |
| m           | $(0 x_{12} 0)$    | 3m1         | $(0 -x_{21} 0)$   |
|             | $(x_{21} 0 x_{23})$ |             | $(x_{21} 0 0)$    |
|             | $(0 0 x_{32})$    |             | $(0 0 x_{33})$    |
| 222         | $(x_{11} 0 0)$    | 6           | $(x_{11} -x_{21})$|
|             | $(0 x_{22} 0)$    |             | $(x_{21} 0 0)$    |
|             | $(0 0 x_{33})$    |             | $(0 0 x_{33})$    |
| mm2         | $(0 x_{12} 0)$    | $-6$        | $(0 0 0)$         |
|             | $(x_{21} 0 0)$    |             | $(0 0 0)$         |
|             | $(0 0 0)$         |             | $(0 0 0)$         |
| 4           | $(x_{11} -x_{21} 0)$ | 622         | $(x_{11} 0 0)$    |
|             | $(x_{11} x_{33})$ |             | $(0 0 x_{33})$    |
|             | $(0 0 0)$         |             | $(0 0 0)$         |
| 222         | $(x_{11} x_{21})$ | 6mm         | $(0 -x_{21} 0)$   |
|             | $(x_{21} -x_{11})$ |             | $(x_{21} 0 0)$    |
|             | $(0 0 0)$         |             | $(0 0 0)$         |
| 422         | $(x_{11} 0 0)$    | $-6m2$      | $(0 0 0)$         |
|             | $(0 x_{11} 0)$    |             | $(0 0 0)$         |
|             | $(0 0 x_{33})$    |             | $(0 0 0)$         |
| 4mm         | $(0 -x_{21} 0)$   | 23          | $(x_{11} 0 0)$    |
|             | $(x_{21} 0 0)$    |             | $(0 0 x_{11})$    |
|             | $(0 0 0)$         |             | $(0 0 x_{11})$    |
| 42m         | $(x_{11} 0 0)$    | 432         | $(x_{11} 0 0)$    |
|             | $(0 -x_{11} 0)$   |             | $(0 0 x_{11})$    |
|             | $(0 0 0)$         |             | $(0 0 x_{11})$    |
| 3           | $(x_{11} -x_{21} 0)$ | $-43m$      | $(0 0 0)$         |
|             | $(x_{21} x_{11})$ |             | $(0 0 0)$         |
|             | $(0 0 x_{33})$    |             | $(0 0 0)$         |

$x_1$ denotes a component of the response tensor $\chi$. The tensor for the triclinic group 1 is not shown, because its form is completely arbitrary. The tensors are given in Cartesian coordinate systems, see the Supplementary Information for details on how these coordinate systems are chosen for each point group. The tensors are given in conventional coordinate systems defined in the International Tables for Crystallography.97

Measurements of spin–orbit fields in NiMnSb

To deduce the vector of the current-induced effective field in NiMnSb, we perform current-driven FMR measurements on two-terminal micro-bars, similarly to previous investigations in spin-valves,52 in ferromagnet/paramagnet bilayers,53, and to our previous experiments in (Ga,Mn)As.10,11

Our samples consist of an insulating InP substrate, 200 nm of an In0.53Ga0.47As buffer-lattice-matched to the substrate and 37 nm of a fully strained NiMnSb film capped with 5 nm of MgO. The material was grown in a multi-chamber molecular-beam-epitaxy system, allowing transfer between different chambers under ultrahigh vacuum.15 The crystal quality of the epilayers was confirmed by high-resolution X-ray diffraction and reflection high-energy electron diffraction (RHEED) measurements. 4 μm × 40 μm bars are patterned by electron beam lithography and ion milling.

The measurement set-up is shown in Fig. 2a. All measurements presented in this work were performed at room temperature using an electromagnet to generate the magnetic field $H_0$, and a rotating stage to set its direction $\phi$ with respect to the bar. The magnetization of NiMnSb is initially aligned along $H_0$. When an alternating electrical current $i(t) = Ie^{i\omega t}$ is passed through the uniformly magnetized micro-bar, carriers gain a non-equilibrium spin polarization due to the ISGE. This spin polarization is exchange-coupled to the magnetization and acts as an effective microwave magnetic field, exciting dynamics.

At resonance, the precessing motion of the magnetization induces a time-dependent change of the bar longitudinal resistance owing to anisotropic magneto-resistance (AMR). Figure 2b shows static measurements of the AMR in bars patterned along [110], [110], [100] and [010] crystal directions. We define $R_{\text{AMR}} = (R(\phi) - R)/R$, where $R$ is the longitudinal resistance and $R$ is the longitudinal resistance averaged over $\phi$. For the [110]- and [100]-oriented bars, $R_{\text{AMR}} \approx C \cos(2\phi)$, where $C \approx 0.06\%$. At resonance, the resistance oscillates with amplitude $dR_{\text{AMR}} \approx -2C \sin(2\phi)\phi$, where $d\phi$ is the precessional amplitude, directly proportional to the torque in the limit of small oscillations.

The product between the oscillating resistance and the oscillating current in the bar yields a rectified component of the longitudinal voltage, $V_{dc} = (dR_{\text{AMR}}(t) \times i(t))$, where $[\times]$ refers to a time average, which is measured via a bias tee. By solving the Landau–Lifshitz–Gilbert equation for an arbitrarily oriented current-induced field, we find that at resonance the rectified voltage is the linear combination of symmetric and antisymmetric Lorentzians, which depend on the components of the field as $^{30,31}$:

$$V_{dc} = JCA_{\text{sym}} \sin(2\phi)h_i$$

$$V_{dc} = JCA_{\text{sym}} \sin(2\phi)[h_i \sin(\phi) + h_c \cos(\phi)]$$

$A_{\text{sym}}$ and $A_{\text{asym}}$ are constants determined by the magnetic anisotropy, whereas $h_i$, $h_c$ and $h_c$ are the components of the current-induced field in the coordinate system specified in Fig. 2a. We note that the $\sin(2\phi)$ term derives from the angle dependence of $dR_{\text{AMR}}$, whereas the $\sin(\phi)$ and $\cos(\phi)$ terms in $V_{dc}^{\text{sym}}$ express the angle dependence of the torque.

Only bars oriented along the [110] and [110] axes were used in the FMR experiments. The AMR in [100]- and [010]-oriented bars nearly vanishes, as seen in Fig. 2b, and cannot produce rectification. In the Supplementary Information we explain, based on symmetry analysis of the AMR measurements and on ab initio AMR calculations, that this is a consequence of the cancellation of non-crystalline and crystalline AMR terms in the NiMnSb film. Nevertheless, FMR measurements in the [110]- and [100]-oriented bars are sufficient for inferring the magnitude and symmetry of the driving spin–orbit fields.

Figure 2c shows $V_{dc}$ for a bar along the [110] crystal direction as the external magnetic field is swept through the FMR condition at different angles to the bar. The resonance is well fitted by an antisymmetric Lorentzian. The independence of the line shape on the frequency of the current indicates that the phase between the current passed in the bar, and the current-induced driving field is fixed and is not affected by reactive components of the circuit.29 This is a necessary requirement to carry the line-shape analysis outlined in refs 10,11. The amplitude of the resonance is proportional to the incident microwave power (see Fig. 2d), implying that the driving field is linear in current density, as is typical for most mechanisms of current-induced torques, including the Oersted field torque and the spin torques.

Figure 3a shows the complete angle analysis of the rectified voltage for a bar along the [110] crystal axis. In these measurements the frequency of the current is fixed at $\omega = 2\pi \times 9$ GHz, with a...
source power of 20 dBm. A resonance is clearly visible at fields above the saturation field of 30 mT. The resonance field depends on the anisotropy of the bar, and its angle dependence can be fitted with the modified Kittel’s formula deduced from the expression of the free energy (for details see Supplementary Information), which at magnetic fields larger than the saturation field reads:

\[
\frac{\omega}{\gamma} = H_1 H_2
\]

where:

\[
H_1 = \mu_0 [H_{\text{res}} + H_{11} \cos \theta - H_{12} \sin 2\theta]
\]

\[
H_2 = \mu_0 \left[ H_{\text{res}} + \frac{H_{11}}{4} (3 + \cos 4\theta) + H_{12}(1 - \sin 2\theta) + M_{\text{eff}} \right]
\]

\(\gamma\) is the gyromagnetic ratio, \(\mu_0 H_{\text{res}}\) is the resonance field, \(\mu_0 H_{11}\) and \(\mu_0 H_{12}\) are the in-plane uniaxial and biaxial anisotropy fields, \(M_{\text{eff}} = M_s - \mu_0 H_{2\perp}\) is the effective magnetization, \(M_s\) being the saturation magnetization and \(\mu_0 H_1\) the out-of-plane anisotropy field, and \(\theta\) is the angle of the external magnetic field.

By rearranging equation (2), we find an expression for \(H_{\text{res}}\) as a function of \(\phi\). Figure 3b,c shows the fitting obtained for the two crystal directions by using this expression. From the fitting we extract the values of the anisotropy fields, treated as free parameters: \(\mu_0 H_{11}^{100} = 638 \pm 3\) mT and \(\mu_0 H_{11}^{110} = 640 \pm 3\) mT, \(\mu_0 H_{12}^{110} = 21.4 \pm 0.3\) mT and \(\mu_0 H_{12}^{100} = 0.1 \pm 0.4\) mT, and \(\mu_0 H_{12}^{010} = 11.1 \pm 0.5\) mT and \(\mu_0 H_{12}^{001} = 7.2 \pm 0.5\) mT. The out-of-plane uniaxial anisotropy field induced by the lattice-matching growth strain is the leading term. Note that the in-plane anisotropies are different in the two micro-bars because their values in the unpatterned film are modified differently by partial (10%) strain relaxation in the 4-μm-wide bars (see Supplementary Information).

In Fig. 4a we plot the amplitude of the resonance with respect to the angle \(\phi\) for the [110]- and [110]-oriented bars. In both cases the amplitude exhibits a \(\sin(2\phi) \cos(\phi)\) dependence, indicating maximum amplitude of precession, thus maximum torque, when the external magnetic field is aligned with the bar. By fitting these graphs with the expression (1) for \(V_{dc}\), we determine the value of the current-induced fields: \(h_{110} = (340 \pm 20)\) μT and \(h_{110} = (550 \pm 50)\) μT normalized to a current density \(J = 10^7\) A cm\(^{-2}\) (the current density was deduced by heating calibration measurements, as detailed in the Supplementary Information).

The current-induced fields measured for the two bars have opposite sign, excluding the Oersted field as a possible driving mechanism of precession and confirming their crystal origin. Figure 4b shows the complete polar plot of the current-induced
field in our bars, constructed by using the measured values $h_{[\bar{1}10]}$ (magenta squares) and $h_{[110]}$ (black circles) as an orthogonal basis. The polar plot is well described by a predominant Dresselhaus field $h_D = (445 \pm 54) \mu$T and a smaller Rashba field $h_R = (105 \pm 54) \mu$T. Similarly to (Ga,Mn)As we attribute the Rashba contribution to symmetry breaking in the crystal structure, which accounts for the uniaxial term of the magnetic anisotropy and is effectively modelled by shear strain in the unpatterned films, combined with partial strain relaxation in the patterned bars.

Measurements on multiple sets of samples patterned along the [110] and [1¯10] crystal directions provided reproducible evidence for a room-temperature field-like torque driven by an effective field with the leading Dresselhaus symmetry.

**Microscopic calculations of spin–orbit fields in NiMnSb**

To obtain a theoretical estimate of the magnitude of the effective field which drives the spin–orbit torque in our NiMnSb film we performed relativistic density functional theory (DFT) calculations using two complementary approaches. In one method, we base our calculations on the full-potential linearized augmented plane-wave (FLAPW) code FLEUR for the description of the electronic structure. The spin–orbit torque is then calculated using the Kubo formalism. In another method, we base our calculations on the self-consistent linearized augmented plane-wave (FLAPW) code and the tight-binding linear muffin-tin orbital approach. In the torque calculation, a scattering region is constructed with the desired disorder and connected to semi-infinite perfectly crystalline leads. The non-equilibrium spin polarization that is carried by conduction electrons is obtained from the explicit scattering wavefunctions using the wavefunction-matching scheme. Thermal disorder is treated in the frozen-phonon approximation. More details on these ab initio methods are given in the Supplementary Information.
In both calculations we considered the symmetry-lowering mechanism of the bulk cubic lattice of NiMnSb due to the substrate-matching growth strain. The resulting current-induced fields are of the Dresselhaus symmetry, in agreement with the crystallographic point-group analysis and with experiment. The magnitude of the current-induced field obtained by the two ab initio methods is 89 and 111 μT per J = 10^4 A cm⁻², respectively. Without any adjustable parameter, the theoretical results agree on the order of magnitude level with our experiments.

**Methods**

Methods, including statements of data availability and any associated accession codes and references, are available in the online version of this paper.

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**Author contributions**

Theory and data modelling were performed by J.G., J.Z., L.Š., Z.Y., J.S., E.F. and T.J. Materials were prepared by F.G. and C.G. Sample preparation was performed by C.C. Experiments and data analysis were carried out by C.G., L.A., V.T. and A.J.F. The manuscript was written by T.J. and C.C., project planning was performed by A.J.F., L.W.M., J.S. and T.J. All authors discussed the results and commented on the paper.

**Additional information**

Supplementary information is available in the online version of the paper. Reprints and permissions information is available online at www.nature.com/reprints. Correspondence and requests for materials should be addressed to A.J.F.

**Competing financial interests**

The authors declare no competing financial interests.
**Methods**

**Materials.** The 37-nm-thick NiMnSb epilayer (room-temperature conductivity of $2 \times 10^4 \, \Omega^{-1} \, \text{cm}^{-1}$) was grown on a 200 nm In$_{0.53}$Ga$_{0.47}$As buffer layer (room-temperature conductivity of 0.1 $\Omega^{-1} \, \text{cm}^{-1}$) on an Fe:InP insulating substrate and capped with 5 nm of MgO. The vertical lattice constant of 5.951 Å indicates a slightly Ni-rich composition.

**Experimental procedure.** A pulse-modulated (at 880 Hz) microwave signal was launched onto a printed circuit board patterned with a coplanar waveguide and then injected into the sample via a bond wire. The rectified voltage, generated at FMR, was separated from the microwave circuit by means of a bias tee, amplified with a voltage amplifier, and then detected with a lock-in amplifier. All measurements were performed at room temperature. A rotating stage allowed setting the orientation of the bar with respect to the fixed in-plane magnetic field generated by an electromagnet.

**Symmetry of the spin–orbit torque.** The spin–orbit torque is caused by a non-equilibrium spin polarization induced by electrical current. This spin polarization interacts with the magnetic moments via an exchange interaction and thus acts on them as a torque. We describe the torque $T$ via an effective field, that is, a field such that

$$ T = M \times B $$

where $B$ is the effective field and $M$ is the magnetic moment.

We describe the non-equilibrium spin polarization using a linear-response theory

$$ \delta s = \chi \cdot E $$

where $\chi$ is a $3 \times 3$ tensor. Similarly we can also write

$$ B = \chi E $$

The non-equilibrium spin polarization and the effective magnetic field have the same symmetry properties. We discuss primarily the field, but all the results apply both for the effective field and the non-equilibrium spin polarization. We thus do not differentiate between $\chi^r$ and $\chi^\perp$, and denote both by $\chi$.

The effective field is usually decomposed into a part even in magnetization and a part odd in magnetization, because these parts have different symmetry properties, affect magnetization dynamics differently, and depend differently on disorder. We discuss only the even part, as it is dominant in the NiMnSb spin–orbit torque experiments as well as in the calculations.

The dependence of the effective field on the magnetization is often described by an expansion in powers of the magnetization. Usually the lowest-order terms in the expansion dominate. Of particular interest is a case when the lowest-order term is independent of magnetization. A torque resulting from such a field is then called field-like, as it induces the same dynamics as an external magnetic field. Although complete analysis of the symmetry of the effective field requires the full magnetic point group, symmetry of the field-like term can be deduced from the non-magnetic point group. If this group contains inversion there can be no non-equilibrium spin polarization or effective field, and therefore also no torque. Table 1 shows field-like terms allowed by symmetry for all point groups with broken inversion symmetry. To find these tensors we used Neumann’s principle—that is, we looked for axial tensors which satisfy

$$ \chi = \det(R)R\chi R^{-1} $$

for all symmetry operations $R$.

The tensors are given in Cartesian coordinate systems. The conventional coordinate systems given in the crystallographic tables are defined by three vectors $a$, $b$, $c$. For the orthorhombic, tetragonal and cubic groups, the vectors $a$, $b$, $c$ are orthogonal, so the choice of the Cartesian system is clear: $x = a/|a|$, $y = b/|b|$, $z = c/|c|$. For the monoclinic groups we consider the unique axis $b$ setting. Then we choose the Cartesian system such that it is right-handed and $x = a/|a|$, $y = b/|b|$. In this case again, the tensors have the same form in the Cartesian system as in the conventional system. For the hexagonal and trigonal groups we choose the right-handed Cartesian system that satisfies $x = a/|a|$, $y = b/|b|$, $z = c/|c|$. The effective field is defined only in magnetic crystals, as it describes a field acting on magnetic moment. The non-equilibrium spin polarization on the other hand can exist in a non-magnetic crystal as well. In such a crystal the field-like term is the only term allowed, because there can obviously be no dependence on magnetization. Therefore, for such a crystal, Table 1 shows exact symmetry of full non-equilibrium spin polarization. In a magnetic crystal it is just a first term in an expansion in powers of magnetization.

**Data availability.** The data that support the plots within this paper and other findings of this study are available from the corresponding author on request.