First-principle approach to the study of spin relaxation times of excited electrons in metals

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We have developed a first-principle method of calculating the spin–lattice relaxation time of excited electrons in metals based on the GW approach incorporating spin–orbit coupling. We have studied the spin–lattice relaxation time and path in Al, Cu, Au, Nb and Ta. The spin–lattice relaxation time and path in Nb appear to be much less than in Al, Cu, Au. They are particularly small in Ta, in accord with strong spin–orbit coupling.

Comparing our results with experimental data on the impurity- and phonon-induced spin relaxation times and path’s, we find that at the energy about of 0.9 eV, which is typical for excited electrons in spin-valve and magnetic-tunnel transistors, the inelastic scattering accompanied by creation of electron–hole pairs is the dominating mechanism of spin–lattice relaxation.

A basic theory of spin relaxation has been elaborated for excited electrons with energy near the Fermi level. It has been proved that four mechanisms are relevant here: Elliott–Yafet, D’yakonov–Perel, Bir–Aronov–Pikus and mechanism of hyperfine interactions [1]. The most important of them seems to be the Elliott–Yafet mechanism where the relaxation of spin polarization is realized through the interaction with impurities at low temperatures and phonons at higher temperatures, with both interactions mediated by spin–orbit coupling (SO).

One may expect that for hot electrons in metals and narrow-band semiconductors at energy above the thermal excitations also important can be the inelastic spin-flip scattering driven by the screened electron–electron interaction. This follows from the great effect of inelastic scattering on the quasi-particle relaxation times in metals, where the decay rate often has a quadratic scaling with respect to energy [2]. Here we develop the first-principle approach to the evaluations of the quasi-particle spin–lattice relaxation times with the scattering mechanism of such kind. The approach is based on the GW method of the many-body perturbation theory [3–5], but takes into account the spin–orbit coupling.

We start from the equation of motion for the Green function $G(x, x')$ of an excited electron inside a solid [3–5]:

$$\left[i(\partial/\partial t) - h(x) - V_{\text{h}}(r, t)\right] G(x, x') = -\int dx'' dr'' \Sigma(x, x''', x') G(x''', x') = \delta(x-x')\delta(t-t'). \quad (1)$$

Here $h(x)$ is the one-electron kinetic energy, $\Sigma(x, x''')$ is the self-energy and $V_{\text{h}}(r, t)$ is the Hartree electrostatic potential

$$V_{\text{h}}(r, t) = \int dr' V_c(r, r') \langle N |\Psi'(x')\Psi(x')|N\rangle = \int dr' V_c(r, r') \rho(r'), \quad (2)$$

which for a many-body state $|N\rangle$ is determined by the usual Coulomb potential $V_c$ and electron density $\rho$. The no-
tations \((x = r, \sigma)\), where \(r\) is the space- and \(\sigma\) is the spin-variable are used here. In stationary conditions the Fourier transform with respect to time \(t \to t'\) converts Eq. (1) into the form

\[
[\omega - h(x) - V_0(x)] G(x, x', \omega) \quad - \int \text{d}x'' \Sigma(x, x'', \omega) G(x'', x', \omega) = \delta(x, x').
\]

The general approach for finding \(G\) is to express it in the spectral representation

\[
G(x, x', \omega) = \sum_i \frac{\psi_i(x, \omega) \psi_i^*(x', \omega)}{\omega - E_i(\omega)},
\]

where \(\psi_i\) are the functions that satisfy the quasi-particle equation

\[
\left[ E_i(\omega) - h(x) - V_0(x) \right] \psi_i(x, \omega) \quad - \int \text{d}x' \Sigma(x, x', \omega) \psi_i(x', \omega) = 0.
\]

The eigenvalues for this equation, quasi-particle energies, are complex, \(E_i = \text{Re} E_i + i \text{Im} E_i\). It is proved in many-body theory that the time evolution of the Green function at \(\omega = E_i\) follows low \(\exp(-\Gamma_i t)\) where \(\Gamma_i = 2 \text{Im} E_i\). Hence, when the Eq. (5) is solved, one has the decay rate of the excitation \(\Gamma_i\) and the quasi-particle lifetime defined as \(\tau_i = \hbar/\Gamma_i\).

The SO coupling, after the Foldy–Wouthuysen transform, has the form [6]:

\[
\hat{H}_\text{SO} = -\frac{\hbar}{4m_e c^2} \hat{\sigma}(E \times \hat{p}),
\]

where \(\hat{\sigma}\) is the vector of Pauli matrices, \(E = -\nabla V\) is the electric field and \(\hat{p}\) is the momentum vector. In the presence of SO coupling the band structure of a paramagnetic solid with the center of symmetry consists of degenerate states with spin components which can be expressed as

\[
\xi_{\uparrow r} = \rho^\uparrow_{\downarrow r} \uparrow + \rho^\downarrow_{\uparrow r} \downarrow,
\]

\[
\xi_{\downarrow r} = \rho^\downarrow_{\uparrow r} \downarrow - \rho^\uparrow_{\downarrow r} \uparrow.
\]

Here \(\rho_{\uparrow r}, \rho_{\downarrow r}\) are the weights of the pure spin-up and spin-down states for the spin-mixed state with wave vector \(p\). Typically \(\rho_{\uparrow r} \gg \rho_{\downarrow r}\), so we have attributed to the states pseudo-spin values \(\zeta = +\) and \(\zeta = -\). Contrary to the screened electron–electron interaction, the spin–orbit coupling in such a form is local and instantaneous, which facilitates its treatment. Namely, one can join the self-energy with SO coupling thus defining the new self-energy

\[
\Sigma_\text{SO}(x t, x' t') = \Sigma(x t, x' t') + \frac{\hbar}{4m_e c^2} \delta(r - r') \delta(t - t') \times \nabla V(r') [p' \times \sigma].
\]

With such re-defined self-energy one obtains for the Green function exactly the same equation of motion (3), but the quasi-particle equation is modified as follows.

In the presence of SO coupling and with spin variable made explicit the quasi-particle Eq. (5) has the form

\[
h_0(r) \psi^\sigma(r, \omega) + \sum_{r'} \int \text{d}r' \Sigma_\text{SO}(r, r', \omega) \psi^\sigma(r', \omega) = E_\sigma(\omega) \psi^\sigma(r, \omega).
\]

Here we have introduced the two-component quasi-particle wave-function \((\psi^\uparrow(r, \omega), \psi^\downarrow(r, \omega))\) and operator \(h_0 = h + V_0\).

In order to solve this equation for a periodic solid one can employ the approximation \(\psi^\sigma(r, \omega) = \psi^\sigma(r)\) where \(\psi^\sigma(r)\) is the \(\sigma\)-component of the spin-mixed band-state at energy \(E^\sigma\), which is calculated by means of the density functional theory with SO coupling included. It is easy to show that Eq. (10) can be reduced to the form

\[
e_{\sigma}^\text{SO} + \sum_\sigma \langle \psi^\sigma \rvert \Sigma_\text{SO}(\omega) \rvert \psi^\sigma \rangle - \sum_\sigma \langle \psi^\sigma \rvert V^x c^\sigma \rvert \psi^\sigma \rangle
\]

\[
- \sum_\sigma \langle \psi^\sigma \rvert H^\text{SO}_c \rvert \psi^\sigma \rangle = E_\sigma(\omega).
\]

Defining \(\Delta\sigma(\omega) = \Sigma_\text{SO}(\omega) - V^x c^\sigma - H^\text{SO}_c\), where \(V^x c^\sigma\) is the LDA exchange-correlation potential, we come to the equation

\[
E_\sigma(\omega) = e_{\sigma}^\text{SO} + \sum_\sigma \langle \psi^\sigma \rvert \Delta\sigma(\omega) \rvert \psi^\sigma \rangle.
\]

So we have the quasi-particle equation similar to that in Refs. [3, 5], but containing spin summation and subtraction of the spin-diagonal terms of the SO coupling.

For the real and imaginary part of the quasi-particle energy it follows then from Eq. (12)

\[
\text{Re } E_\sigma(\omega) = e_{\sigma}^\text{SO} + \sum_\sigma \langle \psi^\sigma \rvert \text{Re } \Delta\sigma(\omega) \rvert \psi^\sigma \rangle,
\]

\[
\text{Im } E_\sigma(\omega) = \sum_\sigma \langle \psi^\sigma \rvert \text{Im } \Delta\sigma(\omega) \rvert \psi^\sigma \rangle.
\]

Applying to Eq. (14) the “on mass shell” approximation [2], i.e. \(\omega = e_{\sigma}^\text{SO}\), we obtain

\[
\text{Im } E_\sigma = \sum_\sigma \langle \psi^\sigma \rvert \text{Im } \Delta\sigma(e_{\sigma}^\text{SO}) \rvert \psi^\sigma \rangle.
\]

The value \(2 \text{Im } E_\sigma\) is the decay rate of the quasi-particle, and its inverse gives the quasi-particle lifetime.

In the case of disentangled spins the general way to find the self-energy and to solve the quasi-particle equation is to employ the formalism of Hedin’s equations for the self-energy [3, 5, 7]. Since the Hedin’s equations for the self-energy are based on the equation of motion, they are not changed with inclusion of the SO coupling. So one can employ the simplest GW approximation of the Hedin’s approach. But in this case the self-energy and Green function are spin-diagonal, therefore the equation for the self-energy in the frequency representation is reduced to

\[
\Sigma_\text{SO}^\sigma(r, r', \omega) = \frac{i}{2\pi} \int \text{d}\omega G_\text{SO}(r, r', \omega + \omega') W(r, r', \omega').
\]
(Hereafter we omit the second \( \sigma \) in the spin-diagonal self-energy.) Now the spectral representation of the Green function (4) includes summation over the states with both pseudo-spin directions, + and −. Then the expectation value of self-energy for the decay of electron in the state \( \psi_{q} \) with momentum \( q \) and energy \( \epsilon_{q} \) in Eq. (15) is expressed as

\[
\langle \text{Im} \Delta \Sigma_{\sigma}(\epsilon_{q}) \rangle = \sum_{k} \sum_{n'} \text{Im} \Delta \Sigma_{\sigma}(q, n; k - q, n').
\]  
(17)

Here we have defined the value

\[
\text{Im} \Delta \Sigma_{\sigma}(q, n; k - q, n') = - \sum_{\sigma'} \text{Im} W_{\sigma \sigma'}(k, \epsilon_{q} - \epsilon_{k + q'}).
\]

and

\[
\times \langle \psi_{q} \psi_{q}^{\sigma'} | B_{n} \rangle \langle B_{n} \psi_{q}^{\sigma'} \psi_{q} \rangle \Theta(\epsilon_{q} - \epsilon_{k + q'}).
\]  
(18)

which is a contribution of the transition from the primary state \( \psi_{q} \) into a secondary lower-energy unoccupied state \( \psi_{k+q'} \). Here we have employed the functions \( B_{n} \) of a complete basis set; \( W \) in Eq. (18) is the square of screened potential in this basis set.

As usually, one can employ the screened potential in random-phase approximation (RPA). In order to obtain this, first one needs to calculate RPA polarization function \( P^{\sigma} \). It has been shown in Ref. [7] that the RPA polarization function in the case of Green functions dependent on two spin coordinates case is expressed as

\[
P^{\sigma}(r_{t1}, r_{t2}) = -i \sum_{\sigma'} G_{\sigma \sigma'}(r_{t1}, r_{t2}) G_{\sigma' \sigma}(r_{t2}, r_{t1}).
\]  
(19)

Applying to Green function the expression (4), making Fourier transforms and expanding in a complete basis set of the functions \( B_{n} \), we obtain the expression for \( P^{\sigma} \):

\[
P_{1}^{\sigma}(q, \omega) =
\sum_{k} \sum_{k'} \sum_{n} \sum_{n'} \frac{1}{\omega - \epsilon_{k + q'} + i\delta} - \frac{1}{\omega + \epsilon_{k + q'} - \epsilon_{k'} - i\delta}
\times \sum_{\sigma, \sigma'} \langle B_{n} \psi_{k}^{\sigma} | \psi_{k'}^{\sigma'} \rangle \langle \psi_{k'}^{\sigma'} | \psi_{k}^{\sigma} B_{n} \rangle.
\]  
(20)

Having calculated \( P^{\sigma} \), one can evaluate the response function \( R = (1 - V^{\sigma} P^{\sigma})^{-1} P^{\sigma} \), inverse dielectric function \( \varepsilon^{-1} = 1 + VR \), screened potential \( W = \varepsilon^{-1} V \) [4] and finally the quasi-particle lifetime via Eqs. (17), (18).

Although the GW approach yields the self-energy diagonal in spin coordinates, one can apply it in the following way in order to evaluate the time of spin-flip \( \tau_{q}^{sf} \). In accord with the Eq. (17) the \( \sigma \)-component of the decay rate in GW approximation can be represented in the form

\[
\text{Im} \Delta \Sigma_{\sigma} = \text{Im} \Delta \Sigma_{\sigma}^{+} + \text{Im} \Delta \Sigma_{\sigma}^{-}.
\]

where the first term includes the summation over pseudospin-majority secondary states, and the second term contains the summation over pseudospin-minority states. Then follows that the value \( \text{Im} E_{q} \) of Eq. (14) which determines the decay rate of a spin-mixed state has four components

\[
\text{Im} E_{q} = \text{Im} \Delta \Sigma_{\sigma}^{+} + \text{Im} \Delta \Sigma_{\sigma}^{-} + \text{Im} \Delta \Sigma_{\sigma}^{+} + \text{Im} \Delta \Sigma_{\sigma}^{-}.
\]  
(21)

Following the Matthiessen’s rule [8], one may associate the first two terms with non-spin-flip decay processes, and the second two terms with the decay accompanied by the flip of spin. Hence the time of spin flip is determined by the formula

\[
\frac{1}{\tau_{q}^{sf}} = -2 \text{Im} \Delta \Sigma_{\sigma}^{sf}(\epsilon_{q})
\]  
(22)

where

\[
\text{Im} \Delta \Sigma_{\sigma}^{sf}(\epsilon_{q}) = \text{Im} \Delta \Sigma_{\sigma}^{+}(\epsilon_{q}) + \text{Im} \Delta \Sigma_{\sigma}^{-}(\epsilon_{q}).
\]  
(23)

In the following we assume that the primary state \( | q n \rangle \) is the pseudospin-majority state, hence the decay processes associated with spin-flips occur to the pseudospin-minority states. Taking into account that both the primary and secondary states are spin-mixed, as it is stated in Eq. (7), we calculate the spin-flip part of the self-energy by means of the Eqs. (14) and (17), but with summation in Eq. (17) over only pseudospin-minority states and the value

\[
\text{Im} \Delta \Sigma_{\sigma}^{+}(q, n; k - q, n') = \text{Im} \Delta \Sigma_{\sigma}^{-}(q, n; k - q, n')
\]  
(24)

replaced with

\[
\text{Im} \Delta \Sigma_{\sigma}^{sf}(q, n; k - q, n') = \text{Im} \Delta \Sigma_{\sigma}^{+}(q, n; k - q, n')
\]

\[
\times (1 - p_{k + q'}) \text{Im} \Delta \Sigma_{\sigma}^{-}(q, n; k - q, n') + p_{k + q'} \text{Im} \Delta \Sigma_{\sigma}^{-}(q, n; k - q, n').
\]  
(24)

where \( p_{k} = |\rho_{k}|^2 \). In the limit of zero SO coupling the defined in such a way value \( \text{Im} \Delta \Sigma_{\sigma}^{sf} \) is identically zero. For a finite SO coupling the first term of the right-hand side counts the probability of the transition from the spin-majority component of the primary state to the spin-minority component of the secondary state and the second term counts the probability of the transitions from the spin-minority component of the primary state to the spin-majority component of the secondary state.

Since the time \( \tau_{q}^{sf} \) is the time of spin reverse into opposite direction, it is reasonable to define the time of complete loss of spin polarization, i.e. the spin–lattice relaxation time \( \tau_{\sigma}^{sr} \), as one half of the spin-flip time. This definition is in accord with the definition introduced in the phenomenological theory of spin-injection in the ferromagnetic metal-normal metal junctions, see Ref. [1].

We calculate the band structure by means of the LMTO method in tight-binding representation [9] with the minimal set of s, p, d-orbitals. In order to accelerate the calculations, we carry out the process of self-consistency neglecting SO coupling. In the calculations of RPA polarization function, response function and screened potential we neglect the SO interactions. The basis set employed in these calculations, and equally in the following self-energy calculations, is constructed from the \( s \times s, s \times p, s \times d, p \times p, p \times d \) products of the LMTO’s by the prescriptions of Ref. [11]; normally the number of such basis states is
45–50. After the calculations of screened potential, the self-consistent band-structure calculations are performed with SO coupling included, as described in Ref. [10], which provides all the necessary data for the subsequent self-energy calculations, in particular the $\rho_{ep}$ and $1-\rho_{ep}$ values. Finally the obtained from the self-energy calculations inverse electron lifetimes and inverse spin-flip times are averaged over momenta in the full Brillouin zone.

In Figs. 1–3 we show the energy dependence of the electron lifetimes and spin-flip times in Al, Cu, Au, Nb and Ta. We have chosen these metals as representatives having different kinds of electronic band structure. Aluminum, whose spin relaxation time has been studied experimentally in Refs. [12, 13] and theoretically in Ref. [14], is a typical free-electron-like metal with relatively low and smooth density of states. Copper and gold, with the spin-flip of for Cu and Au. Notice logarithmic scale for the time values. Nb and Ta. Hence the

The spin-flip times are much higher than the lifetimes; the ratio of the spin-flip time to the lifetime, i.e. $\tau_{sf}/\tau$, appear to be in the limits $8.2 \pm 1.2$ for Al, $7.8 \pm 3$ for Cu, $6.5 \pm 2$ for Au, $18 \pm 2$ for Nb, $8 \pm 2$ for Ta. Hence the energy relaxation is an essentially faster process that the relaxation of spin. Having lost its excess energy after the lifetime, the electron still keeps the spin polarization propagating with the Fermi velocity. A noteworthy result of the calculations is also the absence of correlation between the momentum-averaged spin-flip time and the SO splitting of the band states in the symmetry points of the Brillouin zone which are generally considered to be indications of the strength of SO coupling. (The similar circumstance is valid also for semiconductors GaAs, GaN, as it was noticed in Ref. [21].) A good example is provided by comparison between Cu and Au. In these metals a large splitting takes place for the 6-fold occupied state $\Gamma_{25}$. Due to the SO coupling this state is split into 4-fold an 2-fold states. Our calculations produce the value of SO splitting between these states equal to 0.18 eV in Cu and 1.2 eV in Au. However, we can conclude from the data in Fig. 2 that the momentum-averaged spin-flip time of the electronic states in Cu is only slightly higher than that in Au. This apparent contradiction is explained by the fact that in Cu

Figure 1 Momentum averaged electron lifetime $\tau$ and spin-flip time $\tau_{sf}$ for Al. Notice logarithmic scale for the time values.

Figure 2 Momentum averaged electron lifetime $\tau$ and spin-flip time $\tau_{sf}$ for Cu and Au.

Figure 3 Momentum averaged electron lifetime $\tau$ and spin-flip time $\tau_{sf}$ for Nb and Ta.
and Au the unoccupied free-electron-like states within 5 eV above Fermi are not spin–orbit split.

One more remarkable fact is the great value of the relative spin-flip time $\tau_{sf}/\tau$ in Nb (~18) in comparison with the rest of metals (6.5–8.2). It is associated with the values of the matrix elements of the SO coupling operator $\chi$. In Cu we have for p-states $\langle p|\chi|p\rangle = 0.01$ Ry and for d-states $\langle d|\chi|d\rangle = 0.007$ Ry. For Nb the value of the matrix element is $\langle d|\chi|d\rangle = 0.004$ Ry, two times lower than for the states of Cu. Because of this, the spin-up–spin-down mixing in the low-energy unoccupied d-states of Nb is less than in Cu which provides higher relative spin-flip time. However, the true value of the spin-flip time (see table below) is in Nb lower than in the free-electron-like metals because of low value of lifetime $\tau$ which takes place due to the high density of d-states above Fermi level. In Ta, in spite of the general similarity between the band structures of Nb and Ta, the $\tau_{sf}/\tau$ value and true value of spin-flip time is much less which occur because of the stronger SO coupling, with $\langle d|\chi|d\rangle = 0.124$ Ry.

The calculations of lifetime and spin-flip time permit to calculate also the spin–lattice relaxation path $\lambda_{sr}$ which we define as the distance passed by a hot electron during its spin relaxation time. Since the spin-flip time $\tau_{sf}$ and spin-relaxation time $\tau_{sr} = \tau_{p}/2$ are markedly longer than the lifetime, we suppose that a hot electron at the initial energy $E$ propagates during the lifetime $\tau$ with the group velocity $v(\epsilon)$ and then, having lost its excitation energy, drifts with the Fermi velocity $v_F$. So we have $\lambda_{sr}(E) = v(\epsilon)\tau(E) + [\tau_{sr}(E) - \tau(E)]v_F(\epsilon)$. We calculate velocities $v(\epsilon)$ along the radial direction in Brillouin zone and average them over all the states with energy $E$. The yielded energy dependencies of the spin–lattice relaxation path’s are shown in Fig. 4. In Table 1 and 2 we also compare our calculated spin–lattice relaxation times and path’s with previous experimental data. Our data refer to the energy 0.3 eV above Fermi level and also to 0.9 eV which is the typical hight of barriers in the spin-valve and magnetic tunnel devices [22].

The main mechanisms of spin relaxation for such electrons are the Elliott–Yafet’s scattering with impurities (the $\tau_{imp}$ values) and phonons (the $\tau_{ph}$ values).

We see from Table 1 that our spin–lattice relaxation times at 0.3 eV, and furthermore at 0.9 eV, are much less than the spin relaxation times associated with Elliott–Yafet’s processes. In particular small are our data for Ta at 0.9 eV, due to the strong spin–orbit coupling and high density of states above the Fermi level. As for the theoretical values of the spin–lattice relaxation path, for Al, Cu and Au at 0.3 eV they are comparable with the experimental data on the impurity-induced and phonon-induced path’s. So for the metals with free-electron-like excited states at low excitation energy one may expect comparable effects of the impurity-induced, phonon-induced and electron-induced scattering on the spin relaxation. From Fig. 4 we see, however, that with the rise of energy our spin–lattice relaxation path reduces by a factor of $4 \div 8$, and at 0.9 eV it is much less than the path’s pertinent to the Elliott–Yafet’s mechanisms.

The spin–lattice relaxation path in Nb and Ta appears to be much shorter than in Al, Cu and Au in spite of lower value of the SO matrix element $\langle d|\chi|d\rangle$. This is associated with the differences both in times and velocities. The lifetime and spin-flip time in Nb, Ta are markedly less due to the much higher density of states above the Fermi level. Besides, the group velocity of d-electrons in Nb, Ta is also about 3 times less than the velocity of p-electrons in Cu

<table>
<thead>
<tr>
<th>Material</th>
<th>$\tau_{imp}$ (ps)</th>
<th>$\tau_{ph}$ (ps)</th>
<th>$\tau_{exp}$ (ps)</th>
</tr>
</thead>
<tbody>
<tr>
<td>aluminum</td>
<td>780 [23]</td>
<td>1300 (170)</td>
<td></td>
</tr>
<tr>
<td>copper</td>
<td>560 [16]</td>
<td>1290 (250)</td>
<td></td>
</tr>
<tr>
<td>gold</td>
<td>–</td>
<td>980 (230)</td>
<td></td>
</tr>
<tr>
<td>niobium</td>
<td>780 [17]</td>
<td>160 (30)</td>
<td></td>
</tr>
<tr>
<td>tantalum</td>
<td>–</td>
<td>86 (18)</td>
<td></td>
</tr>
</tbody>
</table>

Table 2 Comparison of experimental spin–lattice relaxation times with our theoretical data. In square brackets the references to the experimental works are given. Our data without brackets refer to the energy of 0.3 eV, and in parenthesis to the energy of 0.9 eV.
and Au. In particular short is the spin–lattice relaxation path in tantalum, which is associated with small velocity, short lifetime and spin-flip time. According to the calculations, the polarization of spin in Ta is lost at the distance about 60 atomic layers. Since this number of layers is typical for spin-valve and magnetic tunnel transistors, tantalum can hardly be effectively used in such devices. Unfortunately, experimental data on spin relaxation in tantalum are absent.

Finally, we have developed a first-principle method of calculating the spin–lattice relaxation times in metals and have applied it to Al, Cu, Au, Nb and Ta. We find that the spin-flip time and spin–lattice relaxation path of the excited free-like electrons in Al, Cu and Au are much longer than those in Nb, Ta. Since the short spin–lattice relaxation path in transition metals is explained not only by the shorter spin-flip time, but also by low velocity and high density of d-states, we expect that any transition metals after Nb with high density of d-state in the vicinity of Fermi level has relatively low spin–lattice relaxation time and path. The calculations also reveal an essential reduction of the spin–lattice relaxation time and path with the energy of excited electrons. The comparison of our results with existing data on the scattering associated with Elliott–Yafet’s mechanism shows that at the energy about 1 eV the inelastic scattering driven by the screened electron–electron interaction is a dominant mechanism of the spin–lattice relaxation.

References