Ab initio magneto-optical Kerr spectra for solid systems with reduction dimensions

A. VERNES*
Center for Computational Materials Science, Technical University Vienna, Gumpendorferstr. 1a, A-1060 Vienna, Austria

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The complex optical conductivity tensor for layered systems is calculated on the basis of the Luttinger formula within the spin-polarized relativistic screened Korringa–Kohn–Rostoker method by means of a contour integration technique. Ab initio Kerr spectra are then obtained for arbitrary geometry and incidence via a $2 \times 2$ matrix technique which includes all multiple reflections between layers and optical interferences in the layers. Applied to fcc Co/Pt layered systems it is shown that the Kerr spectra are highly sensitive to the surface orientation, its dependence on the Co slab and cap Pt thickness closely reflects, e.g. the ABCABC... stacking of the fcc (111) texture, and even for a moderate number of Co/Pt$_3$ bilayers a satisfactory good agreement with corresponding measurements is achieved.

1. Introduction

The recorded signal on a CD, DVD, or magneto-optical (MO) disk is read out by using the magneto-optical Kerr effect (MOKE); i.e. all read-out devices, in principle, exploit that remarkable feature of MOKE. After reflection from oppositely and in magnitude almost equally magnetized domains, the main polarized plane of a linearly polarized incident light is oppositely rotated by almost the same relatively small angle, called Kerr rotation angle [1].

The Co/Pt layered systems due to their perpendicular anisotropy and large Kerr rotation angle in the ultraviolet range of short wavelengths [2], are still seen as promising MO media alternative to the nowadays widely used high-density TbFeCo disks [3]. In spite of the large amount of experimental work done so far (for a review see Carcia and Suzuki [4]), only a few aspects of the mechanisms providing the exceptional MO properties of Co/Pt multilayers have been theoretically addressed. The aim of the present article is to provide further insights to these mechanisms.

This paper is organized as follows. In section 2 the basic theoretical concepts used to determine the inter- and intra-layer contributions to the

*Email: av@cms.tuwien.ac.at
complex optical conductivity within the spin-polarized relativistic screened K\textsuperscript{orringa–Kohn–Rostoker (SKKR) method are briefly described. Section 3 then gives some details concerning the applied numerical schemes, which allows one to fully control the accuracy of the computation. The calculation of Kerr spectra by including all multiple reflections and optical interferences within a layered system via the so-called $2\times2$ matrix technique is presented in section 4. In the case of fcc Co/Pt, the dependence of Kerr angles on the surface orientation (section 5.1), Co slab and cap Pt thickness, respectively (sections 5.2 and 5.3) and also on the repetition rate of Co/Pt\textsubscript{3} bilayers (section 5.4) is discussed in section 5. Finally in section 6 the main results are summarized.

2. Optical conductivity tensor

By using the vector potential description of the electric field, the complex optical conductivity $\tilde{\sigma}_{\mu\nu}\left(\omega\right)$ is given by the Luttinger formula [5],

$$\tilde{\sigma}_{\mu\nu}\left(\omega\right) = \frac{\tilde{\Sigma}_{\mu\nu}(\omega) - \tilde{\Sigma}_{\mu\nu}(0)}{\hbar \omega + i\delta},$$

in terms of the current–current correlation function [6],

$$\tilde{\Sigma}_{\mu\nu}(\omega) = \frac{i\hbar}{V} \sum_{m,n} \frac{f(\epsilon_m) - f(\epsilon_n)}{\epsilon_m - \epsilon_n + \hbar \omega + i\delta} J_{\mu m} J_{\nu n},$$

with $f(\epsilon)$ being the Fermi–Dirac distribution function, $\epsilon_m$ and $\epsilon_n$ a pair of eigenvalues of the one-electron Hamiltonian, $J_{\mu m}$ matrix elements of the electronic current operator ($\mu = x, y, z$), and $V$ the reference (crystalline) volume.

The Luttinger formula (1) has the advantage of simultaneously providing both absorptive and dispersive parts of the optical conductivity tensor on the same footing without using the Kramers–K\text{rö}nig relations [6]. Beside this, the positive infinitesimal $\delta > 0$ introduced to describe the interaction of the system with its surroundings, which makes the fields to be turned on at $t = -\infty$ [7], also accounts for all scattering processes at $T \neq 0$ that are commonly not included in a standard bandstructure calculation [6].

The current–current correlation function $\tilde{\Sigma}_{\mu\nu}(\omega)$ as given by equation (2) is evaluated by performing contour integrations in the complex energy plane at finite temperature [8]. In the selection of contours one exploits the properties of the Fermi–Dirac distribution function $f(\epsilon)$ of complex argument $z \in \mathbb{C}$ [9]. Namely, that $f(\epsilon)$ for $z \in \mathbb{C}$ is analytical everywhere except at the so-called Matsubara poles: $\epsilon_k = \epsilon_F + i(2k - 1)\delta_T$, ($k = 0, \pm 1, \pm 2, \ldots$; $\delta_T = \pi k_B T$) [10], and for energies parallel to the real axis situated in-between the two successive Matsubara poles, $f(\epsilon \pm 2i k \delta_T) = f(\epsilon)$ [11].

Two contours $\Gamma_1$ and $\Gamma_2$ are initially considered around the eigenvalues $\epsilon_m$ and $\epsilon_n$ including a finite number of Matsubara poles. The contour parts of $\Gamma_1$ parallel to the real axis ($z = \epsilon \pm i \delta_j$, $j = 1, 2$) are taken in-between the two successive Matsubara poles $(2N_j - 1)\delta_T < \delta_j < (2N_j + 1)\delta_T$, with $N_j$ being the Matsubara poles in the upper
and \( N_2 \) in the lower semi-plane included in \( \Gamma_1 \), i.e. \( \delta_j = 2N_j\delta_T \) for \( j = 1, 2 \). \( \Gamma_2 \) is \( \Gamma_1 \) mirrored along the real axis and hence includes \( N_2 \) Matsubara poles in the upper and \( N_1 \) in the lower semi-plane. In order not to have \( \epsilon_n - \zeta \) and \( \epsilon_m + \zeta \) included in the contour \( \Gamma_2 \) and \( \Gamma_1 \), respectively, the only constraint to be applied is \( \delta > \delta_2 \).

By using the residue theorem and exploiting the symmetry of contours \( \Gamma_1 \) and \( \Gamma_2 \) [8],

\[
\tilde{\sum}_{\mu\nu}(\omega) = \oint_{\Gamma_1} dz f(z) \tilde{\sum}_{\mu\nu}(z + i\delta, z) - \left[ \oint_{\Gamma_1} dz f(z) \tilde{\sum}_{\mu\nu}(z - \hbar\omega + i\delta, z) \right] ^* \\
- 2i\delta_T \sum_{k=-N_2+1}^{N_1} \left[ \tilde{\sum}_{\mu\nu}(z_k + \hbar\omega + i\delta, z_k) + \tilde{\sum}_{\mu\nu}(z_k - \hbar\omega + i\delta, z_k) \right] 
\]

(3)

and [6]

\[
\tilde{\sum}_{\mu\nu}(0) = \oint_{\Gamma_1} dz f(z) \tilde{\sum}_{\mu\nu}(z, z) - 2i\delta_T \sum_{k=-N_2+1}^{N_1} \tilde{\sum}_{\mu\nu}(z_k, z_k), 
\]

(4)

where in terms of the electronic Green function \( G(z) \), the kernel \( \tilde{\Sigma}_{\mu\nu}(z_1, z_2) \) is written

\[
\tilde{\sum}_{\mu\nu}(z_1, z_2) = -\frac{\hbar}{2\pi v^*} \text{Tr}[J^\mu G(z_1) J^\nu G(z_2)].
\]

(5)

Due to the trace entering this expression, the optical conductivity tensor, as given by equation (1) and calculated by means of equations (3) and (4), includes both inter- and intra-band contributions on the same footing in a parameter-free manner [6].

In the following, the computation of the inter- and intra-layer contributions \( \tilde{\Sigma}_{pq}(z_1, z_2) \) via equation (5) is performed by using relativistic current operators [12] and the Green function as obtained within the spin-polarized relativistic SKKR method for layered systems [13]. Hence by using equation (1) one ends up with this purely quantum mechanical part of the calculations by obtaining the inter- and intra-layer contributions \( \tilde{\sigma}_{pq}(\omega) \) to the conductivity tensor. In the case of plane waves it has been shown [14] that from these \( \tilde{\sigma}_{pq}(\omega) \) quantities the corresponding contributions to the permittivity are directly obtained, e.g. within the Gaussian system of units, as

\[
\varepsilon_{pq}(\omega) = \delta_{pq} I + \frac{4\pi i}{\omega} \tilde{\sigma}_{pq}(\omega) \quad \text{for} \quad p, q = 1, \ldots, N,
\]

(6)

with \( \delta_{pq} \) being the Kronecker symbol, \( I \) the \( 3 \times 3 \) identity matrix, \( \tilde{\omega} = \omega - i\delta \) the complex frequency, and \( N \) denoting the number of layers in the system.

3. Computational details

From numerical point of view, besides of Matsubara poles, \( \tilde{\Sigma}_{pq}(\omega) \) depends also on the number of complex energy points \( n_z \) considered for energy integrals in equations (3) and (4), and on the number of \( k \)-points used to calculate the scattering path operator within SKKR and \( \tilde{\Sigma}_{pq}(z \pm \hbar\omega + i\delta, z) \) for a given energy \( z \), respectively. For this reason, two efficient schemes to control the accuracy of these \( z \)- and \( \bar{k} \)-integration have been introduced [15].
The first one is controlling the accuracy of \( z \)-integrations along each contour part by comparing the results obtained using the Konrod quadrature [16, 17], \( \mathcal{K}_{2n+1} \sum_{\mu \nu}^{\text{pq}}(\omega) \), with those provided by the Gauss integration rule, \( G_n \sum_{\mu \nu}^{\text{pq}}(\omega) \) [18]. On a particular part of the contour, \( \sum_{\mu \nu}^{\text{pq}}(\omega) \) is considered to be converged, if the criterion [15]:

\[
\max \left| \mathcal{K}_{2n+1} \sum_{\mu \nu}^{\text{pq}}(\omega) - G_n \sum_{\mu \nu}^{\text{pq}}(\omega) \right| \leq \varepsilon_z,
\]

is fulfilled for a given \( \varepsilon_z \) defining the accuracy.

The other scheme permits the two-dimensional \( \vec{k} \)-space integrations to be performed with an arbitrary high precision \( \varepsilon_{\vec{k}} \) within the cumulative special points method [15]. The applied convergence criterion:

\[
\max \left| \mathcal{S}_{n_1} \sum_{\mu \nu}^{\text{pq}}(z', z) - \mathcal{S}_{n_{i-1}} \sum_{\mu \nu}^{\text{pq}}(z', z) \right| \leq \varepsilon_{\vec{k}},
\]

is fulfilled during the computation for each complex energy \( z \) on the contour or \( z_k \) Matsubara pole. Note that here \( n_i = 2^{i+2} n_0 (n_0 \in N) \) is the number of divisions along each primitive translation vector in the two-dimensional \( \vec{k} \)-space.

In the present article, all optical conductivity tensor calculations were carried out such as the convergence criteria (7) and (8) were fulfilled for \( \varepsilon_z = \varepsilon_{\vec{k}} = 10^{-3} \) a.u. Due to the independence of \( \sum_{\mu \nu}^{\text{pq}}(\omega) \) of the contour path in the upper semi-plane [15], on the other hand, the computational time can be significantly reduced by imposing \( N_1 \gg N_2 \) for \( N_2 \) as small as possible.

Not discussed before, is the \( T \)-dependence of the optical conductivity tensor when it is obtained from the Luttinger formula (1) by using equations (3) and (4) for the current–current correlation functions. In fact, \( \delta_T \) providing the imaginary part of the Matsubara poles is the only quantity in the present scheme, which directly depends on the electronic temperature \( T \) involved by the Fermi–Dirac distribution. Considering \( \delta = 2\delta_2 \) in accordance with the only constraint applied in selecting the contours, recall section 2, if \( T \) is varied, the life-time broadening \( \delta \) too is changing. However, in order to preserve the broadening of bands involved in optical transitions, by keeping \( \delta \) of 0.048 Ryd (0.653 eV) [15], the considered number of Matsubara poles in the lower semi-plane \( N_2 \) has to be \( T \)-dependent, see table 1. Furthermore, as mentioned earlier, to accelerate the computation one has to have the difference

<table>
<thead>
<tr>
<th>( T ) (K)</th>
<th>( \delta_T = \pi k_B T ) (mRyd)</th>
<th>( N_2 )</th>
<th>( \delta = 4N_2 \delta_T ) (Ryd)</th>
<th>( N_1 = N_2 + 35 )</th>
</tr>
</thead>
<tbody>
<tr>
<td>50</td>
<td>0.9948</td>
<td>12</td>
<td>0.048</td>
<td>47</td>
</tr>
<tr>
<td>85</td>
<td>1.6912</td>
<td>7</td>
<td>0.047</td>
<td>42</td>
</tr>
<tr>
<td>100</td>
<td>1.9897</td>
<td>6</td>
<td>0.048</td>
<td>41</td>
</tr>
<tr>
<td>120</td>
<td>2.3876</td>
<td>5</td>
<td>0.048</td>
<td>40</td>
</tr>
<tr>
<td>150</td>
<td>2.9545</td>
<td>4</td>
<td>0.048</td>
<td>39</td>
</tr>
<tr>
<td>200</td>
<td>3.9794</td>
<td>3</td>
<td>0.048</td>
<td>38</td>
</tr>
<tr>
<td>300</td>
<td>5.9693</td>
<td>2</td>
<td>0.048</td>
<td>37</td>
</tr>
</tbody>
</table>
$N_1 - N_2$ large enough, e.g. 35 poles, and hence besides of $N_2$, also $N_1$ must be adapted to the magnitude of $T$.

Calculations of the optical conductivity tensor performed for the parameters given in table 1 in the case of the fcc Co/Pt/Pt (100) have proven that excepting some low photon energy regions, the $T$-dependence of the layer-resolved permittivities as defined in equations (6) and (19) can be neglected. For the sake of simplicity, in figure 1, only the difference in the off-diagonal tensor element $\tilde{\varepsilon}_{xy}$ with respect to $T = 300$ K data is depicted for the surface Co layer ($p = 10$). As it can be seen from this figure, the $T$-dependence of $\tilde{\varepsilon}_{xy}$ is showing up only for photon energies below 2 eV. Similarly, in the case of the diagonal tensor element $\tilde{\varepsilon}_{xx}$, not shown here, a weak $T$-dependence occurs when $\omega < 1$ eV. Although $\tilde{\varepsilon}_{xx}$ for the surface Co and the

![Figure 1](image-url)

**Figure 1.** Difference in the off-diagonal layer-resolved permittivity tensor element $\Delta \tilde{\varepsilon}_{xy} = \tilde{\varepsilon}_{xy}(T_m) - \tilde{\varepsilon}_{xy}(T)$, where $T_m = 300$ K, for the surface Co layer ($p = 10$) in the case of fcc Co/Pt/Pt(100) layered system as a function of the photon energy $\omega$ and electronic temperature $T$. 
Pt layer immediately below this \((p = 9)\) are almost identical in shape, only the real part of \(\varepsilon^{p}_{xx}\) is of the same magnitude. The off-diagonal tensor element \(\varepsilon^{p}_{xy}\), on the other hand, completely differs for \(p = 9\) and 10 and that \(\varepsilon^{p}_{xy} \neq 0\) for \(p = 9\), uniquely proves the polarizability of the fcc Pt(100) substrate. Interesting, but not surprising, is that the high polarizability of the Pt – at most five atomic monolayers – as demonstrated by Vernes et al. [19], is not altered at all by \(T\). For the paramagnetic semi-infinite fcc Pt(100) substrate below the polarized Pt buffer layers, the permittivity \(\varepsilon\) is almost independent of \(T\) for all photon energies \(\omega\) within the visible range. But relatively large differences in comparison with the \(T = 300\) K data are observed for low photon energies (below 1 eV for \(\text{Re} \varepsilon\) and \(\omega < 2\) eV for \(\text{Im} \varepsilon\), see figure 2. These differences, however, range to an extent, which does not exceed the magnitude of differences shown by the diagonal tensor elements of the layer-resolved permittivities in the case of Co/Pt\(_9\)/Pt(100).

Figure 2. Difference in the permittivity \(\Delta \varepsilon = \varepsilon(T_m) - \varepsilon(T)\), where \(T_m = 300\) K, in the case of fcc Pt(100) substrate as a function of the photon energy \(\omega\) and electronic temperature \(T\).
4. 2 × 2 Matrix technique

From the optical point of view each layer \( p \) is assumed to be a homogeneous, linear, and anisotropic conducting medium characterized by the complex permittivity \( \varepsilon^p(\omega) \). In the case of cubic, hexagonal, or tetragonal systems, for example, if the orientation of the magnetization \( \vec{M}_p \) points along the surface normal (z-direction), the layer-resolved permittivity tensor is approximately written [20] as

\[
\varepsilon^p(\omega) = \begin{pmatrix}
\varepsilon_{xx}^p(\omega) & \varepsilon_{xy}^p(\omega) & 0 \\
-\varepsilon_{xy}^p(\omega) & \varepsilon_{yy}^p(\omega) & 0 \\
0 & 0 & \varepsilon_{zz}^p(\omega)
\end{pmatrix},
\]

(9)

by supposing that \( \varepsilon_{zz}^p(\omega) \approx \varepsilon_{xx}^p(\omega) \). Note that the error introduced by this assumption is directly proportional to \( \varepsilon_{zz}^p(\omega) - \varepsilon_{xx}^p(\omega) \), which in turn is small enough to be neglected. In fact, calculations in the case of fcc Co/Pt5/Pt(111) for parameters as in the last entry of table 1 have shown that independently on the particular layer index \( p \), the higher the \( \omega \), the smaller the \( \varepsilon_{zz}^p(\omega) - \varepsilon_{xx}^p(\omega) \) and the relative error \( |\varepsilon_{zz}^p(\omega) - \varepsilon_{xx}^p(\omega)|/\varepsilon_{xx}^p(\omega) \) does not exceed 10% even for relatively low photon energies.

In the following, the layers are numbered starting with the first one on top of the semi-infinite substrate \((p = 1)\), such that if \( N \) layers are considered, the index of the surface layer is \( p = N \). It is also convenient to label the substrate and the vacuum by 0 and \( N + 1 \), respectively.

With the permittivity \( \varepsilon^p(\omega) \) as given by equation (9), the normal modes for the propagation of the electric and magnetic plane waves in a layer \( p \) are directly obtained by solving the Fresnel (characteristic) equation [21], which is of fourth order in the \( z \) Cartesian component \( \tilde{n}_{pz} \) of the complex refraction. Note that the other two Cartesian components [22],

\[
\begin{align*}
\tilde{n}_{px} &= -\sin\theta \cos\varphi = \tilde{n}_x, \\
\tilde{n}_{py} &= -\sin\theta \sin\varphi = \tilde{n}_y
\end{align*}
\]

(10)

are all known in terms of the spherical colatitude \( \theta \) \((0 \leq \theta < \pi/2)\) and longitude \( \varphi \) \((0 \leq \varphi \leq 2\pi)\) of the incident light coming in from the vacuum side. For each normal mode \( \tilde{n}_{pq}^{(k)} \) \((k = 1, \ldots, 4)\), the Helmholtz equation [23] then immediately provides all Cartesian components of the electric field \( \vec{E}_p^{(k)} \), which in turn substituted into the curl Maxwell equation [24], finally yield the magnetic field \( \vec{H}_p^{(k)} \).

In practice, however, the determination of \( \vec{E}_p^{(k)} \) is complicated by the fact, that the Helmholtz equation for a given \( \tilde{n}_{pq}^{(k)} \) has to be solved by keeping at least one Cartesian component of the electric field arbitrary. Among all possible parameterization of the electric fields, a physically very transparent scheme results by following Mansuripur’s strategy [25, 26]:

\[
\begin{align*}
\varepsilon_{px}^{(k)} &= \text{arbitrary} \\
\varepsilon_{py}^{(k)} &= \alpha_p^{(k)} \varepsilon_{px}^{(k)}, \quad \text{for } k = 1, 3 \\
\varepsilon_{pz}^{(k)} &= \beta_p^{(k)} \varepsilon_{px}^{(k)} \\
\varepsilon_{px}^{(k)} &= \alpha_p^{(k)} \varepsilon_{py}^{(k)} \\
\varepsilon_{py}^{(k)} &= \text{arbitrary}, \quad \text{for } k = 2, 4.
\end{align*}
\]

(11)
By exploiting in each layer the continuity of the tangential components of the total electric and magnetic field at the lower boundary \( z_p, \ p = 0, \ldots, (N + 1) \), the layer-resolved reflectivity matrix \( R_p \), which relates all arbitrary electric field components to each other [27],

\[
R_p = (D_{p-1}A_p^{34} - B_p^{34})^{-1}(B_p^{12} - D_{p-1}A_p^{12}),
\]

(12)
can be determined recursively by starting from the vanishing reflectivity matrix \( R_0 = 0 \) of the substrate and by means of the following \( 2 \times 2 \) matrices:

\[
A_{p}^{k,k+1} = \begin{pmatrix} 1 & \alpha_p^{(k+1)} \\ \alpha_p^{(k)} & 1 \end{pmatrix}, \quad B_{p}^{k,k+1} = \begin{pmatrix} \alpha_p^{(k)} - \beta_p^{(k)} & \alpha_p^{(k+1)} - \beta_p^{(k+1)} \\ -\beta_p^{(k)} & \alpha_p^{(k)} - \beta_p^{(k)} \end{pmatrix}.
\]

(13)

Furthermore, also the propagation matrices,

\[
c_{p}^{k,k+1} = \begin{pmatrix} \exp\left[+i\theta_p \alpha_p^{(k+1)} d_p\right] & 0 \\ 0 & \exp\left[+i\theta_p \alpha_p^{(k+1)} d_p\right] \end{pmatrix}, \quad (k = 1, 3 \text{ and } p = 0, \ldots, N),
\]

(14)

with \( d_p = z_{p+1} - z_p > 0 \) being the thickness of layer \( p \), are needed to construct the auxiliary matrices:

\[
D_p = (B_p^{12}c_p^{12} + B_p^{34}c_p^{34}R_p)(A_p^{12}c_p^{12} + A_p^{34}c_p^{34}R_p)^{-1}, \quad p = 0, \ldots, N.
\]

(15)

In vacuum there are only two normal modes, namely an incident and a reflected beam, therefore the surface reflectivity matrix is given by

\[
R_{\text{surf}} = R_{N+1} = (D_N - B_{N+1}^{34})^{-1}(B_{N+1}^{12} - D_N) = \begin{pmatrix} \tilde{r}_{xx} & \tilde{r}_{xy} \\ \tilde{r}_{yx} & \tilde{r}_{yy} \end{pmatrix},
\]

(16)

where in terms of equation (13),

\[
B_{N+1}^{12} = -B_{N+1}^{34} = \frac{1}{\sqrt{1 - (\tilde{n}_x)^2 + (\tilde{n}_y)^2}} \begin{pmatrix} \tilde{n}_x\tilde{n}_y & 1 - (\tilde{n}_x)^2 \\ -1 + (\tilde{n}_y)^2 & -\tilde{n}_x\tilde{n}_y \end{pmatrix}.
\]

(17)

Because \( R_{\text{surf}} \) relates the tangential components of the reflection light to the corresponding components of the incident light, the Kerr angles \( \theta_K \) and \( \epsilon_K \), respectively, in terms of the complex reflectivity coefficients \( \tilde{r}_{\mu\nu} \), e.g. for a \( p \)-polarized incident light (\( \theta \) arbitrary and \( \varphi = \pi/2 \)) are simply given by

\[
\theta_K + i\epsilon_K \simeq -\frac{\tilde{r}_{xy}}{\tilde{r}_{yy}} \cos \theta,
\]

(18)

for more details see Vernes and Weinberger [27].

By using the \( T \)- and \( \omega \)-dependent layer-resolved permittivities calculated for each entry of table 1 in the case of fcc Pt(100) substrate (\( p = 0 \)) and Co/Pt\(_9\) layered
system \((p=1, \ldots, 10)\), the aforesaid \(2 \times 2\) matrix technique immediately yields \(T\)-dependent polar Kerr spectra. For normal incidence, i.e. for \(\theta=0\), differences of these spectra with respect to those obtained for \(T=300\) \(K\) are shown in figure 3. Obviously, in the low photon energy region, where all layer-resolved permittivities \((p=0, \ldots, 10)\) are quite sensitive to the selection of \(T\) (see figures 1 and 2), the polar Kerr angles are also \(T\)-dependent. However, in contrast to layer-resolved permittivities of which diagonal tensor elements for low \(\omega\) are much stronger influenced by \(T\) than the off-diagonal tensor elements, the \(T\)-dependence of the polar Kerr spectra for normal incidence can be en masse neglected for all photon energies. First, because the changes in the polar Kerr spectra with \(T\) are at least one order of magnitude smaller than the Kerr angles themselves, for example, in the case of Co/Pt\(_9\)/Pt(100)

![Graph](image-url)

Figure 3. Differences in the polar Kerr rotation \(\Delta \theta_K = \theta_K(T) - \theta_K(T_m)\) and ellipticity \(\Delta \varepsilon_K = \varepsilon_K(T) - \varepsilon_K(T_m)\) angles for normal incidence, where \(T_m=300\) \(K\), in the case of fcc Co/Pt\(_9\)/Pt(100) as a function of the photon energy \(\omega\) and electronic temperature \(T\). Here 1 mdeg = 10\(^{-3}\) deg.
these differences do not exceed on average 0.003 deg, see figure 3. Second, because the only effect of $T$ on the Kerr spectra is a shift of angles by an almost $\omega$ independent small amount, which hence preserves all main characteristics of the spectra, like its shape, position of peaks, and the zero axis crossing points.

Using the above technique together with linear material equations, it has been shown [20] that the layer-resolved permittivities $\tilde{\varepsilon}^p(\omega)$ can be self-consistently obtained – at least for polar geometry, when equation (9) holds – from the contributions $\tilde{\varepsilon}^{pq}(\omega)$ as given by equation (6). Since at least in the case of polar geometry and normal incidence [20], and for longitudinal geometry and oblique incidence [22], very little differences between the zeroth order for the layer-resolved permittivities,

$$\tilde{\varepsilon}^p(\omega) = \sum_{q=1}^{N} \tilde{\varepsilon}^{pq}(\omega) \quad \text{for} \quad p = 1, \ldots, N$$

and their self-consistent counterparts arise, in the following instead of using a self-consistent procedure simply equations (6) and (19) are used.

The advantage of using the zeroth–order permittivity over the self-consistent one, compare figure 4 with the corresponding figure 3 from Vernes et al. [14], is manifold already in the case of fcc Pt bulk. (For both calculations the parameters from the last entry of table 1 were used.) As it can be seen from figure 4, due to the lower values below 2 eV of the zeroth–order permittivity with respect to those which were self-consistently determined, $\text{Re} \tilde{\varepsilon}$ agrees better with the experimental data. In the case of $\text{Im} \tilde{\varepsilon}$, below 2 eV the zeroth–order data are unfortunately much lower than the experimental ones and the sharp double peak structure seen in experiments around 0.7 eV is theoretically only reproduced by the (100) oriented surface. But the strong divergence of $\tilde{\varepsilon}$ obtained self-consistently for photon frequencies $\omega < \delta = 0.653$ eV, within the zeroth–order approach is completely removed.

5. Results and discussions

All results presented in this section concern the parameterization provided by the last entry of table 1. Ab initio Kerr spectra were calculated for different fcc Co/Pt layered systems such that the total number $N$ of atomic layers considered was always kept at a multiple of three. Note that this rather technical requirement is a direct consequence of the special form of the structure constants involved by the SKKR method [13].

Considering that the semi-infinite fcc Pt bulk (lattice parameter of 3.92 Å) acts as a charge reservoir for the layered system on its top, the self-consistently determined Fermi level of Pt bulk (approximately of $-0.03$ Ryd slightly depending on the surface orientation) has been fixed for the layered system too and used in all calculations.

5.1. Surface orientation

Kerr spectra of a single Co atomic monolayer (ML) on top of semi-infinite fcc Pt(100), Pt(110) and Pt(111) bulk, respectively, have been calculated including
Figure 4. Zeroth–order symmetrized permittivity of fcc Pt bulk for the (100), (110), and (111) surface orientation (diamonds, circles, and squares) as obtained for $T = 300 \text{ K}$ and using the Pt/Pt$_{12}$/Pt layered system. The experimental data are taken from Weaver [28] – crosses, Hunter et al. [29] – pluses, Seignac and Robin [30] – stars, Kirillova et al. [31] – right triangles and Yu et al. [32] – down triangles.
five Pt buffer layers. As it can be seen from figure 5, the polar Kerr angles for normal incidence are highly sensitive on the orientation of the surface. Although the frequency dependence of the Kerr rotation angle for the (110) surface orientation $\theta_{K}^{(110)}(\omega)$ is similar to that obtained for the (100) surface orientation $\theta_{K}^{(100)}(\omega)$, below 3 eV: $\theta_{K}^{(110)} < \theta_{K}^{(100)}$, while the opposite holds beyond 3 eV. The local maximum around 5 eV of the Kerr rotation angle is for $\theta_{K}^{(110)}(\omega)$ in photon energy higher positioned $\theta_{K}^{(110)}(\omega)$ than the corresponding maximum of $\theta_{K}^{(100)}(\omega)$, but for both cases of surface orientation a local minimum occurs at exactly 3 eV.

The Kerr rotation angle for the (111) surface orientation $\theta_{K}^{(111)}(\omega)$, on the other hand, shows a completely different $\omega$-dependence than $\theta_{K}^{(110)}(\omega)$ or $\theta_{K}^{(100)}(\omega)$. Although $\theta_{K}^{(111)}(\omega)$ too has a local minimum at 3 eV, in the case of the (111) surface orientation, two maxima of approximately equal high exist around 3.5 and 5.5 eV. Because the self-consistently calculated layer-resolved magnetic moments, not shown here, are almost identical for the (100) and (111) surface oriented Co/Pt layered systems, one immediately concludes that the peculiarities of $\theta_{K}^{(111)}(\omega)$ are mainly due to its (surface) electronic properties imposed by the (111) texture.

Excepting the sharpness of the negative peak around 4.5 eV, $\epsilon_{K}^{(100)}(\omega)$ and $\epsilon_{K}^{(110)}(\omega)$ as a function of the photon energy are very similar beyond 1 eV, but $\epsilon_{K}^{(110)}(\omega) > \epsilon_{K}^{(100)}(\omega)$. Furthermore, below 1 eV and $\epsilon_{K}^{(110)}(\omega)$ are almost indistinguishable.

The Kerr ellipticity angle for the (111) texture, $\epsilon_{K}^{(111)}(\omega)$, is smaller than $\epsilon_{K}^{(100)}(\omega)$ below 1 eV, in-between 3 and 3.5 eV, and in the range of photon energies starting from 4.5 eV up to the end of visible regime at 6 eV. Although $\epsilon_{K}^{(111)}(\omega)$ too achieves its maximum around 1 eV, where also the Kerr ellipticity angles for the other two surface orientations do, no global minimum of $\epsilon_{K}^{(111)}(\omega)$ exists around 5 eV. In fact, beyond 1 eV, the Kerr ellipticity angle in the case of the (111) surface orientation continuously decreases in the visible regime (also slightly oscillating above 3 eV).

The calculated Kerr angles for the (111) surface orientation are in a qualitatively good agreement with the experimental data from van Drent and Suzuki [33, 34]. Although the peaks in the theoretical spectra are all centered at lower photon energies than in experiments, the measured oscillations of Kerr angles in the case of Co(3 Å)/Pt(11 Å) layers are well reproduced by the calculations. A reason of this shift in energy is the Pt substrate used in calculations, which has a significantly smaller lattice constant (3.92 Å) than Si (5.43 Å) or 80 Å SiN used as support for the layered samples during the deposition. Thus quite obviously the bandstructure of all Co/Pt layered systems investigated in here is wider than that of the measured probes [33]. Quantitatively, the calculated Kerr angles are approximatively 15 times smaller in magnitude than the experimental ones. This relatively huge difference is not surprising at all, if one admits that the Kerr rotation angle should linearly increase with the magnetization and recalls the dimension of samples used in the experiments. Indeed, for a (111) surface orientation, i.e. for a perpendicular stacking thickness of 2.26 Å, the Co(3 Å)/Pt(x Å) samples of 250 Å thickness correspond to (Co$_{1,32}$/Pt$_{x}$)$_{110,37}$ with $x = 1.32$, 2.21, and 4.86 ML, of which magnetization quite clearly has to be larger than that of Co/Pt$_{5}$.

The dip around 3.7 eV – reproduced by the calculations with a shift of 0.5 eV towards the lower photon energies – is ascribed by van Drent and Suzuki [34] to the
Figure 5. Polar Kerr angles for normal incidence in the case of fcc Co/Pt for (100), (110), and (111) surface orientations (open diamonds, circles, and grey squares). The properly scaled ($\times 1/15$) experimental data for Co(3 Å)/Pt($x$ Å) with a total thickness of 250 Å and $x = 3$ (full circles), 5 (full squares), and 11 (full diamonds), respectively, are taken from van Drent and Suzuki [33, 34].
high polarization of the Pt layers below the ferromagnetic Co surface. But the same authors could not uniquely assign the other negative peak in the ultra-violet range around 5 eV to the contribution arising from the Co [33, 34]. No doubt, that the substrate contributes substantially to the Kerr spectra. However, the occurrence of a similar local minimum in the calculated Kerr rotation angle, makes one to conclude that this is rather originated on the magneto-optical activity of Co (recall the Kerr spectra of fcc Co – not shown here), than arising due to the Pt substrate.

Another parameter on which the Kerr angles depend, is the angle of incidence \( \theta \), which is especially important, because no Kerr instrumentation exists, where a perfect normal incidence could be worked out. The Kerr measurements for Co/Pt as quoted above, for example, were performed for an angle of incidence \( \theta = 4^\circ \), see figure 1 of van Drent and Suzuki [33]. Recently, it has been proven [27] that the smallest possible departure from a perfect polar geometry and normal incidence can enormously enhance the sensitivity of Kerr angles on the polarization state of the incident light. Namely, for a non-vanishing angle of incidence, the smallest deviation of the incident light from the linear polarization will show up in the Kerr spectra too. For this reason, in figure 6 the \( p \)-polarization and the spherical longitude \( \phi = 90^\circ \)

![Figure 6](image_url)

Figure 6. Polar Kerr angles for oblique incidence of a \( p \)-polarized light in the case of fcc Co/Pt(111).
of the incident light are kept fixed, while the angle of incidence, i.e. the spherical colatitude \( \theta \), is continuously changed from a perfect normal incidence up to 90°. In all these calculations, the surface has been considered (111) oriented, because it has been experimentally demonstrated that Pt as substrate promotes this particular fcc(111) texture [35]. Note that for figure 5 it does not have to be specified which linearly polarized incident light has been used, because in the ideal case of polar geometry and normal incidence, obviously, the \( s \)- and \( p \)-waves yield identical Kerr spectra [27].

As it can be seen in figure 6, independently on the photon energy considered within the visible range, in magnitude the largest Kerr rotation angle is obtained for \( \theta = 57° \) and the largest Kerr ellipticity angle for an incidence angle of 75°, respectively. Furthermore, one observes that by increasing the angle of incidence to 57°, the decrease of \( \theta^{(111)}(\omega) \) is less pronounced in the low energy region \( (\omega \leq 3 \text{ eV}) \) than in the high one of \( \omega \geq 3 \text{ eV} \). Another interesting aspect is that the dip, in the ultraviolet 4.5 eV with increasing \( \theta \), is moving to 5 eV and is simultaneously getting deeper, exactly as seen in the measurements [33, 34]. For incidence angles beyond 57°, the Kerr rotation angle independently on \( \omega \) is approaching the zero axis and completely vanishes for \( \theta = 90° \). The Kerr ellipticity angle \( \epsilon^{(111)}_K(\omega) \), on the other hand, without changing its shape over the whole visible range of frequencies is monotonously increasing for an increasing angle of incidence below 75°, while beyond 75° it is continuously decreasing to zero, which value is reached for \( \theta = 90° \).

5.2. Co slab on top of fcc Pt(111)

Commonly, the Kerr rotation angle for normal incidence as a function of the Co slab thickness on the top of fcc Pt(111) substrate has two main features: a negative peak around 3.5 eV and a sharp decay at the end of the visible regime \( (\sim 6 \text{ eV}) \), see figure 7, which show up in the spectrum, when one has at least four Co surface layers. The local maximum in-between these two dips at about 4.5 eV touches the zero axis for 8 and 11 Co atomic monolayers and becomes positive valued for more than 11 Co ML. In fact, the Kerr rotation spectra beyond 12 Co surface layers are all very similar, almost identical.

In the case of the Kerr ellipticity angle for normal incidence the two dominant positive peaks, around 3 and 5.5 eV in the visible spectrum of photon energies, occur when there are at least six Co layers on the Pt substrate and with the increasing number of Co layers both increase in magnitude.

Unfortunately, there are no available experimental data known in the literature for fcc Co\(_{N'}\)/Pt(111), \( N' = 1, \ldots, 16 \text{ ML} \). Compared with the experimental polar Kerr spectra of fcc Co, not shown here, one immediately realizes that none of the peaks characterizing the calculated spectra can be uniquely ascribed to the contribution of the Co slab alone. Which could mean that even with a further increase of the Co slab thickness on top of Pt(111), the magneto-optical properties of fcc Co could not be recovered.

Depending on the photon energy, the polar Kerr angles for normal incidence as a function of the number of Co layers \( N' \) on top of Pt(111) are oscillating with an \( N' \)-dependent amplitude and a period of 2–4 Co ML around either a constant or with \( N' \) quadratically changing average value. The oscillations of the Kerr
rotation angle $\theta_K(\omega; N')$ for $\omega < 2.72$ eV, in particular, can be neglected, while on average $\theta_K(\omega; N')$ for $2.72 \leq \omega \leq 4.08$ eV and $\omega \geq 5.44$ eV decreases (in absolute value increases) and for $4.08 < \omega < 5.44$ eV increases (in magnitude increases) once the Co slab thickness is increased. In the case of the Kerr ellipticity angle $\epsilon_K(\omega; N')$, on the other hand, independently on the photon energy $\omega$, its average always increases by increasing $N'$.

By comparing figure 7 with figure 8, it directly follows that the polar Kerr angles for normal incidence reproduce well the 3 Co ML oscillations of the magnetization in fcc Co$_N$/Pt(111) layered systems. This means that the generalized Fresnel coefficient [36], which relates the Kerr rotation angle to the magnitude of the magnetization, is almost independent of the Co slab thickness and hence the oscillations of the Kerr angles are predominantly of magnetic origin closely reflecting the ABCABC... stacking of the fcc (111) surface. Another interesting feature of Co$_N$/Pt(111) layered systems is that independent of the number of Co ML, only the
same few (approximately three) Pt layers below the Co slab are significantly polarized. Note that depending on the Co slab thickness, the total number of Pt buffer layers is either 5, 6, or 7, such that the first Pt buffer layer on top of the Pt substrate has always a vanishing magnetic moment – with an accuracy of $\pm 0.01 \mu_B$.

In order to figure out how these oscillations of the polar Kerr angles depend on the angle of incidence a photon energy of approximately $5.986507 \pm 6$ eV, for which the largest oscillations occur at normal incidence – see figure 7, has been considered. The calculated polar Kerr angles for oblique incidence and this photon energy are given in figure 9. By inspecting this figure, one immediately observes that no significant differences occur for the polar Kerr rotation angle with respect to that obtained for normal incidence, if the angle of incidence $\theta$ remains below $45^\circ$. By increasing the angle of incidence beyond $45^\circ$, e.g. for $\theta = 50^\circ$ the polar Kerr rotation angle decreases, if $N' \leq 7$ Co ML and increases when $N' > 7$. A further increase of the incidence angle mainly preserves this tendency, but the higher $\theta$ the lower the limit for the Co thickness from starting which hugely enlarges the Kerr rotation angle against the polar Kerr rotation angle for normal incidence. Excepting the case of two Co atomic monolayers on top of fcc Pt(111), i.e. for all $N' \neq 2$, the polar Kerr rotation angle is maximized when $\theta \approx 80^\circ$. Furthermore, oscillations
of the polar Kerr rotation angle as a function of $N$ seen for $\theta < 45^\circ$ persist for higher angles of incidence too, but they become smeared out in some extent.

The dependence of the polar Kerr ellipticity angle on the angle of incidence is somewhat simpler than that of the Kerr rotation angle: by continuously increasing $\theta$ below $60^\circ$, the polar Kerr ellipticity angle is always shifted to higher values, which just slightly depends on $N'$. As long as $\theta \leq 60^\circ$, the small positive peak seen for $N' = 4$ at normal incidence completely disappears and a new local maximum occurs for the layered system containing 12 Co ML. This latter peak persists even for angles of incidence beyond $60^\circ$, for which the polar Kerr ellipticity angle decreases continuously till $\theta$ equals $90^\circ$, when quite obviously the Kerr effect vanishes.

5.3. Pt cap layers on top of Co/Pt (111)

Experimentally it has been shown that Pt cap layers on the top Co layer deposited on fcc Pt(111) are needed to prevent the oxidation of the surface [37]. As shown in figure 10, the smallest amount of Pt, say already 1 ML, on top of Co surface layer
dramatically changes the polar Kerr spectra for normal incidence with respect to the
corresponding uncapped layered system. In particular, the presence of Pt cap layers
significantly lowers (in magnitude increases) the Kerr rotation angle over the entire
visible range of photon frequencies. Up to three Pt cap layers the shape of the Kerr
rotation spectrum as known for the uncapped Co/Pt(111), see figure 5, is more or less
preserved. But three large dips are showing up in the Kerr rotation spectrum around
2, 4, and 5.5 eV, if there are four to six Pt cap layers on top of Co surface. For more
than six Pt layers, these negative peaks are getting less pronounced and the first
one, i.e. around 2 eV, is moved towards higher photon energies such that starting
with seven and up to twelve Pt cap layers, the Kerr rotation angle as a function
of the photon energy $\omega$ – after the strong decay in-between 0 and approximately
2 eV – is almost perfectly cosine-like.

The changes due to the Pt cap layers in the polar Kerr ellipticity angle for normal
incidence with respect to the uncapped Co/Pt(111) system are less dramatic than for
the Kerr rotation angle. Excepting two positive peaks around 3.5 and 5 eV, which
occur for one up to six Pt cap layers, the polar Kerr ellipticity spectra possess the

Figure 10. Polar Kerr angles for normal incidence in the case of Pt capped fcc Co/Pt(111).
same characteristics (e.g. zero axis crossing point) as the uncapped one. These positive peaks almost completely disappear from the Kerr ellipticity spectra once one has at least seven Pt cap layers on top of Co. In fact, beyond the zero axis crossing point around 2 eV, the Kerr ellipticity angle decreases almost linearly for seven up to twelve Pt cap layers.

Independently on the photon energy $\omega$, the polar Kerr rotation angle for normal incidence as a function of Pt cap layers oscillates with a period of approximately 3 ML typical for the ABCABC... stacking of the fcc (111) surface. The amplitude of these oscillations is strongly $\omega$ – and number of cap layers dependent. Excepting a small region of about 3 eV, for all other photon energies a global minimum of the polar Kerr rotation angle for normal incidence is achieved by using four Pt cap layers. In the middle of the visible regime (i.e. around 3 eV) this global minimum is provided by the layered system having two Pt cap layers on top of Co. In contrast to the Kerr rotation angle, the polar Kerr ellipticity angle for normal incidence as a function of the Pt cap layers, for more than four cap layers irregularly oscillates in both amplitude and period around an almost constant average value, which in turn strongly depends on the photon energy. For less than four Pt cap layers, a sharp dip occurs for 3 ML in the polar Kerr ellipticity spectra, which is then moved to 2 ML for photon energies at the end of the visible regime.

5.4. Co/Pt$_3$ bilayers

As it can be seen in figure 11, by increasing the number of Co/Pt$_3$ bilayers, the polar Kerr rotation angle for normal incidence continuously and almost linearly decreases, i.e. in magnitude increases, with a slope directly proportional to the photon energy. The negative peak at the end of the visible regime of photon frequencies around 5.5 eV present already in the Kerr rotation spectrum of fcc Pt$_3$/(Co/Pt$_3$)/Pt(111) is getting deeper and deeper by including more and more bilayers in the system. In addition, another dip at about 4.25 eV is showing up once the layered system contains more than five Co/Pt$_3$ bilayers. In Kerr measurements, this ultraviolet negative peak is found more pronounced and broader than the calculated one and it moves from 4.1 eV to 3.9 eV in these experiments by increasing the thickness of the Co sublayer [38, 39]. Further measurement have shown that this negative peak around 4 eV – less pronounced within the Kerr rotation spectrum of fcc Co, but representative for Co/Pt layered systems – occurs mainly due to the interaction between Co and Pt [40]. As previously proved [19], the Pt masks the magneto-optical activity of Co within the Co/Pt layered systems and therefore it is not surprising, that the typical infra-red sharp dip at about 1.5 eV provided by fcc Co – also observed in the case of thick Co films on top of fcc Pt(111), see figure 7 – almost disappears from the polar Kerr rotation spectra of Pt$_3$/(Co/Pt$_3$)$_x$/Pt (111) layered systems.

Although the obtained positive peak around 1 eV agrees well well the experimental one, the most prominent difference between the calculated and measured polar Kerr ellipticity angle for normal incidence is the absence of an other peak at about 3 eV from the theoretical spectra, see figure 11. As a consequence of this, one notes that the zero axis crossing points as seen in the calculations are shifted by approximately 1 eV to lower photon energies in comparison with the measured ones, which are all situated around 3.7 eV [40]. The decay of the polar Kerr ellipticity
Figure 11. Polar Kerr angles for normal incidence in the case of fcc Pt$_3$/(Co/Pt$_3$)$_N$/Pt(111). Circles, squares, diamonds, and triangles (up, left, down, and right) mark the calculated data for $N' = 1, \ldots, 6$ (open symbols) and the measured ones from Uba et al. [38, 39] – grey and full symbols, respectively – for layered systems containing at least 30 Co$_2$/Pt$_6$, Co$_{2.6}$/Pt$_7$, Co$_3$/Pt$_5$, Co$_4$/Pt$_6$, and Co$_8$/Pt$_6$ bilayers.
beyond these points, on the other hand, is stronger for the experimental spectra. However, in contrast to fcc Co, for both measured and calculated spectra the minimum Kerr ellipticity angle is achieved somewhere out of the visible range of photon energies.

Altogether, makes one to conclude that in spite of the moderate number of Co/Pt₃ bilayers $N''=1,\ldots,7$ considered in the calculations, the main features of the polar Kerr angles for normal incidence as seen in the experiments are well reproduced. Note that the repetition rate of bilayers in the Kerr measurements used samples, in a fcc (111) texture corresponds to Co₂/Pt₆, Co₂.₆/Pt₇, Co₃/Pt₅, Co₄/Pt₆, and Co₈/Pt₆, is at least of 33 [39].

Independently on the photon energy, the polar Kerr rotation angle for oblique incidence in figure 12 decreases for an angle of incidence $\theta$ smaller than $55^\circ$ and beyond this angle of incidence increases up to zero achieved for $\theta=90^\circ$. For all angles of incidence between 0 and $90^\circ$, the polar Kerr rotation angle as a function of the number of Co/Pt₃ bilayers linearly decreases with different slopes for $N''=1,\ldots,4$ and $N''=5,\ldots,7$. The same applies also to the polar Kerr ellipticity

![Graph](image-url)

Figure 12. Polar Kerr angles for oblique incidence ($p$-polarized light of $\omega\approx6\,\text{eV}$) in the case of fcc Pt₃/(Co/Pt₃)$_{N''}$/Pt(111).
angle for oblique incidence, expecting the fact that this continuously increases when
the angle of incidence is increased from 0 up to \(75^\circ\), whereas beyond this angle of
incidence it decreases.

6. Summary

Concerning the computation of the complex optical conductivity tensor on the basis
of the Luttinger formula by applying a contour integration technique within the
spin-polarized relativistic screened Korringa–Kohn–Rostoker, it has been proven
that excepting the low photon energy region, the \(T\)-dependence of the
layered–resolved permittivities is negligible. By using then these quantities as
input to calculate the polar Kerr angles for normal incidence via the so-called
2 \(\times\) 2 matrix technique, it was found that Kerr spectra for the entire visible range
of photon energies are \(T\) independent with an accuracy of 0.003 deg.

Having one Co ML on top of differently oriented semi-infinite Pt bulk, it has
been shown that the Kerr spectra for normal incidence are very sensitive on the
surface orientation. In the case of Co/Pt(111), it has been found that independently
on the photon energy, the largest polar Kerr rotation (ellipticity) angle occurs for an
incidence angle of \(57^\circ\) (75\(^\circ\)).

By varying the Co slab thickness on the top of Pt(111) substrate, oscillating polar
Kerr angles were obtained independently on the photon energy or incidence angle, of
which period is in accordance with the ABCABC . . . stacking of the fcc (111) texture.
For a photon energy at the end of the visible regime, independently on the Co slab
thickness, the Kerr rotations angle is maximized for an angle of incidence of
approximately 80\(^\circ\), whereas the Kerr ellipticity angle for 60\(^\circ\).

The Pt cap layers on the surface Co dramatically change the Kerr spectra with
respect to the uncapped Co/Pt(111) layered systems. As a general feature, one notes
that the Kerr rotation angle decreases (in magnitude increases), while the Kerr
ellipticity angle remains almost unchanged, if the Co/Pt(111) is capped by using Pt.

Most closely to the capped Co/Pt(111) samples widely prepared for Kerr
experiments are those containing Co/Pt3 bilayers. Therefore, it is not surprising
that a comparison of the calculated Kerr spectra for these layered systems with
the measured ones, not only shows a satisfactory good agreement between theory
and experiments, but also emphasizes those differences, which occur when the
dimension of the layered systems is significantly reduced against measured probes.
For example, it has been shown that independently on the photon energy used, the
Kerr angles linearly vary with the number of Co/Pt3 bilayers and that this feature is
preserved for all angles of incidence between 0 and 90\(^\circ\). If \(\omega \approx 6\ \text{eV}\), the maximum of
Kerr angles is achieved for an incidence angle of 55\(^\circ\) and 75\(^\circ\) in the case of the Kerr
rotation and ellipticity, respectively.

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