Large magneto-optical Kerr effect in noncollinear antiferromagnets Mn$_3$X ($X =$ Rh, Ir, or Pt)

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(Dated: September 10, 2015)

Magneto-optical Kerr effect, normally found in magnetic materials with nonzero magnetization such as ferromagnets and ferrimagnets, has been known for more than a century. Here, using first-principles density functional theory, we demonstrate large magneto-optical Kerr effect in high temperature noncollinear antiferromagnets Mn$_3$X ($X =$ Rh, Ir, or Pt), in contrast to usual wisdom. The calculated Kerr rotation angles are large, being comparable to that of transition metal magnets such as bcc Fe. The large Kerr rotation angles and ellipitici- ties are found to originate from the lifting of the band double-degeneracy due to the absence of spatial symmetry in the Mn$_3$X noncollinear antiferromagnets which together with the time-reversal symmetry would preserve the Kramers theorem. Our results indicate that Mn$_3$X would provide a rare material platform for exploration of subtle magneto-optical phenomena in noncollinear magnetic materials without net magnetization.

PACS numbers: 71.15.Rf, 75.30.-m, 75.50.Ee, 78.20.Ls

I. INTRODUCTION

Magneto-optical coupling effects reflecting the interactions between light and magnetism are fundamental phenomena in solid state physics. Originally, Faraday$^2$ and Kerr$^3$ discovered, respectively, that when a linearly polarized light beam hits a magnetic material, the polarization plane of the transmitted and reflected light beams rotates. Although magneto-optical Faraday and Kerr effects have been known for over a century, they have become the subjects of intense investigations only in the past decades, mainly due to the applications of optical means in modern high-density data-storage technology. Faraday effect attracts less attention than Kerr effect because it can only occur in ultra-thin films, where complexities of multiple reflections and discontinuous polarizations at the interfaces with the substrate arise. In contrast, magneto-optical Kerr effect (MOKE) is widely used as a powerful probe of the electronic and magnetic properties of materials, such as the domain wall$^{5,6}$, surface plasma resonance$^{7,8}$, magnetic anisotropy$^{9,10}$ and topological insulator$^{11,12}$.

Band exchange splitting caused by magnetization together with relativistic spin-orbit coupling (SOC) has been recognized as the origin of MOKE$^{13,14}$. Therefore, MOKE has been explored extensively in various ferromagnetic transition metals as well as their alloys and compounds. By ferromagnets here we meant the magnetic materials with net magnetization including ferrimagnets. Experimentally, Erskine and Stern$^{15,16}$ first reported the MOKE spectra of bulk Fe, Co, Ni, and Gd, and discussed their relationships with $d$-band widths and electron spin polarizations. After that, large Kerr rotation angles of $\sim 1.0^\circ$ were observed in several Mn-based ferromagnetic alloys, such as PtMnSb$^{22}$, MnBi$^{23}$ and MnPt$_3$.$^{24}$ On the theoretical side, first-principles density functional calculations can directly capture the MOKE spectra with an impressive accuracy compared to experiments. For instance, Guo and Ebert$^{16,17}$ studied the the MOKE spectra in bulk Fe and Co as well as their multilayers. Kim$^{25}$ et al. focused on the surface effect and structural dependence of the MOKE spectra in Co thin films and CoPt alloys, and Stroppa$^{18}$ et al. analyzed the electronic structure and magneto-optical property of uniformly Mn-doped GaAs. Very recently, Rosa$^{19}$ et al. also performed the study of the magneto-optical property of Mn-doped GaAs in a special digital ferromagnetic heterostructure. Moreover, Ravindran$^{20}$ et al. investigated the magnetic, optical, and magneto-optical properties of manganese pnictides and found a systematic increase of the Kerr rotation angles from MnAs to MnSb, and to MnBi.

Although MOKE experiments have been almost always conducted on various types of ferromagnets in the past$^{14,24}$, no explicit conclusion has been established that MOKE must be absent when either magnetization or SOC is not present. In particular, whether MOKE can arise from a spin non-polarized system (without magnetization) or not is still an open question. Remarkably, Chen$^{25}$ et al. recently revealed that the anomalous Hall effect, which has a physical origin akin to that of MOKE, is large in noncollinear antiferromagnet Mn$_3$Ir with zero net magnetization. This surprising result stems from the fact that in a three-sublattice kagome lattice with a noncollinear triangle antiferromagnetic configuration, not only the time-reversal symmetry $T$ is broken but also there is no spatial symmetry operation $S$ which, in combination with $T$, i.e., the $TS$, is a good symmetry that preserves the Kramers theorem. In other words, band exchange splitting exists in this system despite of zero net magnetization. Naturally, it would be interesting to explore possible MOKE in Mn$_3$Ir as well as its isostructural materials Mn$_3$Rh and Mn$_3$Pt, which are widely considered as the promising candidates in information-storage devices due to their prominent exchange-bias properties$^{26}$ and high Néel temperatures$^{27,28}$.

In this paper, we present a comprehensive first-principles
study of MOKE in spin non-polarized systems, focusing on antiferromagnets $\text{Mn}_3X$ ($X = \text{Rh}$, Ir, or Pt) in two low-energy noncollinear spin structures. We show that, because of the strong SOC and the breaking of band double-degeneracy, the absorption rates of the left- and right-circularly polarized lights differ significantly, giving rise to a previously undetected MOKE in these antiferromagnets without net magnetization. The Kerr rotation angles of $\text{Mn}_3X$ increase from $X = \text{Rh}$, to Ir, and to Pt, due to the increased SOC strength of the $X$ atom. The largest one of $\sim 0.6^\circ$ in $\text{Mn}_3\text{Pt}$ is comparable to those of elemental transition metals, such as Fe and Co, reported previously$^{16,17}$. Our first-principles calculations also show that the MOKE would diminish if the SOC is switched off, demonstrating the essential role of the SOC for the occurrence of the MOKE in these systems. Our theoretical results suggest noncollinear antiferromagnets $\text{Mn}_3X$ to be an interesting material platform for further studies of novel magnetooptical phenomena and technological applications.

II. THEORY AND COMPUTATIONAL DETAILS

MOKE generally refers to the change in the polarization property of light when it interacts with magnetism, that is, a linearly polarized light shone on the surface of a magnetic sample will become elliptically polarized in the reflected beam. The ellipticity $\varepsilon_K$ and Kerr angle $\theta_K$ (rotation of the major axis relative to the polarization axis of the incident beam) are widely used to probe and characterize magnetic materials. $\varepsilon_K$ and $\theta_K$ are usually combined into the complex Kerr angle, $\phi_K = \theta_K + i\varepsilon_K$. Depending on the directions of photon propagation and magnetization vector with respect to the surface plane, there are three different geometries of the Kerr effect, namely, the polar, longitudinal, and transverse geometries. Of these, the polar geometry usually has the largest complex Kerr angle and thus is the most interesting one in connection with technological applications. In this paper, we consider the polar geometry as a prototype and the other two geometries can be obtained similarly.

For a solid with at least threefold rotational symmetry, the elements of optical conductivity tensor satisfy $\sigma_{xx} = \sigma_{yy}$ and $\sigma_{xy} = -\sigma_{yx}$. In such a case, the absorptive parts of optical conductivity tensor (real diagonal and imaginary off-diagonal elements) due to interband transitions, can be obtained using the Kubo’s formula within the linear response theory$^{20-22}$

$$\sigma_{xx}^1(\omega) = \frac{\lambda}{\omega} \sum_{k,j,j'} \left[ \Pi_{jj'}^+ \right] \left[ \Pi_{jj'}^- \right] \delta(\omega - \omega_{jj'})$$(1)

$$\sigma_{xy}^2(\omega) = \frac{\lambda}{\omega} \sum_{k,j,j'} \left[ \Pi_{jj'}^+ \right] \left[ -\Pi_{jj'}^- \right] \delta(\omega - \omega_{jj'})$$ (2)

where $\lambda = \frac{\pi e^2}{2\hbar^2 v_F^2}$ is a material specific constant, $\hbar \omega$ the photon energy, $\hbar \omega_{jj'}$ the energy difference between the occupied and unoccupied bands at the same $k$-point, and $\Pi_{jj'}^\pm = \langle \mathbf{k} \mid \left( \frac{\hbar}{\sqrt{2}} p_y \pm i \hbar p_y \right) \mathbf{k} \rangle^j \mathbf{k}^j$ the dipole matrix elements for circularly polarized light with $+$ and $-$ helicity, respectively. The corresponding dispersive parts, namely, the imaginary diagonal element $\sigma_{xx}^2(\omega)$ and real off-diagonal element $\sigma_{xy}^1(\omega)$, can be obtained by use of the Kramer-Kronig transformation$^{23}$

$$\theta_K + i\varepsilon_K = \frac{-\sigma_{xy}}{\sigma_{xx} \sqrt{1 + (4\pi/\omega)^2 \sigma_{xx}}},$$ (3)

which can be explicitly evaluated from the optical conductivity tensor calculated from the electronic structure of the solid concerned. Since the intraband transitions contribute little to the off-diagonal elements of optical conductivity tensor in the magnetically ordered materials$^6$ and affect mainly the MOKE spectra below $1\sim 2$ eV$^{16,17,19}$, here we take into account only the interband transition contribution as expressed in Eqs. (1-3).

In this paper, we consider ordered cubic $L1_2$ $\text{Mn}_3\text{Rh}$, $\text{Mn}_3\text{Ir}$, and $\text{Mn}_3\text{Pt}$ alloys and adopt the experimental lattice constants of $3.813 \text{ Å}^{28}$, $3.785 \text{ Å}^{35}$ and $3.833 \text{ Å}^{22}$, respectively. The total energy and electronic structure are calculated based on first-principles density functional theory with the generalized-gradient approximation in the form of Perdew-Berke-Ernzerhof$^{38}$. The accurate frozen-core full-potential projector-augmented wave method$^{21}$, as implemented in the Vienna $ab$ initio simulation package (VASP)$^{38,39}$, is used. The fully relativistic projector-augmented potentials are adopted in order to include the SOC. The valence configurations of Mn, Rh, Ir, and Pt atoms taken into account in the calculations are $3d^64s^1$, $4d^85s^1$, $5d^86s^1$, and $5d^66s^1$, respectively. A large plane-wave energy cutoff of 350 eV and a fine Monkhorst-Pack $k$-point mesh of $16\times16\times16$ are used for the self-consistent electronic structure calculations. For the calculation of optical conductivity tensors, a denser $k$-point mesh of $20\times20\times26$ (8833 $k$-points) in the irreducible

- Figure 1. (Color online) Cubic $L1_2$ crystal structure of $\text{Mn}_3X$ ($X = \text{Rh}$, Ir, or Pt) with T1 (a) and T2 (b) spin configurations, as well as the corresponding (111) planes in (c) and (d), respectively. The red and green balls represent Mn and $X$ atoms, respectively. The dashed lines in (c) and (d) stand for the primitive cell of the kagome lattice.
Brillouin zone) is used in the tetrahedron integration.

III. RESULTS AND DISCUSSION

In this section, we first present the calculated total energy and magnetic properties of the low energy noncollinear magnetic structures (T1, T2, and T3) of Mn$_3$X ($X$ = Rh, Ir, or Pt) and also compared our results with available previous reports in Sec. III A. Then, the calculated optical conductivity, the key ingredient for evaluating the MOKE, for the two low energy spin structures (T1 and T2) is reported in Sec. III B. Finally, we present the large magneto-optical Ker effect in the ordered Mn$_3$X alloys in Sec. III C.

A. Magnetic structure

There are two kinds of crystal structures for Mn$_3$X alloys, namely, ordered $L1_2$-type and disordered $\gamma$-phase. The ordered alloys were found to be noncollinear antiferromagnetic with one of the two nearly degenerate spin configurations, namely, the T1 and T2 triangle structures, as shown in Figs. 1(a) and 1(b), respectively. The Mn spin magnetic moments basically lie in the (111) plane and point to the center (along the edge) of the triangle, forming three nearest-neighboring Mn sublattices, for the T1 (T2) configuration, which can be viewed as two-dimensional kagome lattices shown in Figs. 1(c) and 1(d). Due to strong exchange-interactions acting on the Mn moments, the Néel temperatures in these Mn-based alloys is as high as 475 K in Mn$_3$Pt, 855 K in Mn$_3$Rh, and 960 K in Mn$_3$Ir.

|          | Energy (meV) | $m_{Mn}$ ($\mu_B$/atom) | $m_{tot}$ ((111)$_B$) |
|----------|--------------|--------------------------|------------------------|
| Mn$_3$Rh | T1           | 3.17, 3.62               | -0.001                 |
|          | T2           | 3.18                     | 0.002                  |
|          | T3           | 3.19, 3.16               | 0.000                  |
| Mn$_3$Ir | T1           | 2.96, 2.94               | -0.029                 |
|          | T2           | 2.96                     | 0.027                  |
|          | T3           | 2.97                     | 0.000                  |
| Mn$_3$Pt | T1           | 3.12, 3.06               | -0.013                 |
|          | T2           | 3.11                     | 0.012                  |
|          | T3           | 3.15, 2.93               | 0.000                  |

TABLE I. The calculated total energies and spin magnetic moments for the T1, T2 and T3 magnetic structures of the Mn$_3$X alloys. Small total spin magnetic moments, being parallel to the (111) direction and due to the spin-canting caused by SOC, exist. The X atoms have a zero magnetic moment, dictated by the site-symmetry of their positions in the crystal structure. The available previously reported Mn spin moments are also listed for comparison. Note that the T2 structure reported in Ref. 40 is named the T3 structure here because it has a higher energy than the T2 structure here. The listed total energies are relative to that of the T1 state.

The calculated total energy and spin magnetic moments together with previously reported experimental and theoretical results are listed in Table I. Clearly, T1 is the magnetic ground state of the Mn$_3$X alloys, being in good agreement with previous experimental and theoretical works. Nevertheless, the energy difference between the T1 and T2 spin configurations is small (being in the order of $\sim$1 meV), i.e., the T1 and T2 are nearly degenerate. Indeed, in the absence of the SOC, all the T1, T2 and T3 spin structures have the same total energy because they are equivalent. Note that the T2 configuration here is not the same as the T2 magnetic structure reported in Ref. 40. Our total energy calculations show that for Mn$_3$Ir, the total energy of the T2 structure in Ref. 40 is $\sim$6.4 meV higher than the T2 configuration here, and thus should be referred to as the T3 spin configuration.

The calculated spin magnetic moment of the X atom is always zero, due to its special site-symmetry, while those of the Mn atoms have nearly identical values of $\sim 3\mu_B$ for all the three states. The calculated spin magnetic moments agree fairly well with previous reports. Further inspecting the total magnetization, we find a nonvanishing component along the $(\langle 111 \rangle \langle 111 \rangle)$ direction for T1 (T2) states, because the Mn moments rotate slightly away from the (111) plane within a very small angle of $\sim 0.1^\circ$. Nonetheless, the Mn$_3$X alloys could still be safely considered as spin non-polarized systems in the sense that the net total magnetization is very small and hardly affects the physical quantities of interest here, such as optical conductivities and MOKE spectra, as will be discussed in the next subsection. Note that this small spin-canting is caused by the presence of the SOC. Interestingly, such a small spin-canting does not occur in the T3 structure (Table I).
sures the corresponding difference. In Figs. 2(a) and 2(b), we
ion of left- and right-circularly polarized light while
of 0.1 eV to simulate the finite lifetime effect of electron.

The adsorptive parts of optical conductivity, i.e., \( \sigma^1_{xx} \) and
\( \sigma^2_{xy} \), have direct physical interpretations. From Eqs. (1) and
(2), it is clear that \( \sigma^1_{xx} \) measures the average in the absorp-
tion of left- and right-circularly polarized light while \( \sigma^2_{xy} \) mea-
sures the corresponding difference. In Figs. 2(a) and 2(b), we
show the \( \sigma^1_{xx} \) for the T1 and T2 spin configurations in the
energy range of 0–6 eV, respectively. Since the \( \sigma^1_{xx} \) is directly
related to the joint density of states and interband transition
probability but does not depend strongly on the details of the
spin structure, it is not surprising that the calculated \( \sigma^1_{xx} \)
are nearly the same for the two different spin configurations.
Moreover, the \( \sigma^1_{xx} \) for all the three Mn\(_3\)X alloys have similar
behaviors, mainly due to their isostuctural nature. In particu-
lar, the \( \sigma^1_{xx} \) for all the three alloys has a prominent broad
peak centered at 2.5 eV. The \( \sigma^2_{xy} \) for the T1 and T2 configura-
tions are displayed in Figs. 2(c) and 2(d), respectively. For
both configurations, the \( \sigma^2_{xy} \) of all the three Mn\(_3\)X alloys have
pronounced oscillatory peaks in the low energy region and its
magnitude reduces gradually to a small value above 6 eV (not
shown). For each individual Mn\(_3\)X alloy, the \( \sigma^2_{xy} \) for the T1
and T2 states differ in sign, due to the opposite total spin mag-
etic moments (see Table I), although they are similar in line
shape and magnitude. Positive (negative) \( \sigma^2_{xy} \) suggests that
the interband transitions are dominated by the excitations due
to the left (right) circularly polarized light. Interestingly, the
sign of the \( \sigma^2_{xy} \) for the T1 structure can be reversed by re-
versing the Mn spin moments while that for the T2 structure
remains unchanged when the chirality of the spin structure is
reversed.

B. Optical conductivity

Physically speaking, the \( dc \) limit of the imaginary off-diagonal element of optical conductivity, \( \sigma^2_{xy} (\omega = 0) \), is
nothing but the anomalous Hall conductivity, which can
also be precisely evaluated by the integration of the Berry
curvature over the Brillouin zone. Chen et al. recently
pointed out that the anomalous Hall effect can arise from
non-collinear antiferromagnet Mn\(_3\)Ir in the T1 spin structure with-
out net magnetization due to the absence of certain spatial
symmetries. This could be understood in terms of the fact that
in the kagome lattice [the (111) plane of the Mn\(_3\)X alloys, as
shown in Figs. (c) and (d)] there is no spatial symmetry \( S \) such as mirror and rotation that in combination with the
time-reversal symmetry \( T \) (i.e., the \( ST \)) can be a good sym-
metry such that the band Kramers degeneracy will be kept in
the system with broken time-reversal symmetry (TRS). This
is certainly in contrast to the case of, e.g., a collinear bipar-
tite antiferromagnet, where the combination of the translational
operation of half of a lattice vector with the time-reversal oper-
ation is a good symmetry that will preserve the band Kramers
degeneracy despite of the broken TRS due to the antiferro-
magnetism. This lifting of the band Kramers degeneracy to-
gether with the strong SOC in the Mn\(_3\)X alloys gives rise to
the nonzero anomalous Hall conductivity. Similarly, one can
expect that the \( \sigma^2_{xy} \) at optical frequencies would be nonzero
as well and from Eq. 3, result in nontrivial magneto-optical
Kerr effect in the Mn\(_3\)X alloys, which will be discussed in
next subsection. Of course, one may argue that the nonzero
\( \sigma^2_{xy} \) could be due to the nonzero total spin magnetic moment
in the T1 and T2 spin structures (Table I). To clarify this, we
also calculate the \( \sigma^2_{xy} \) spectrum from the electronic structure
with a zero spin magnetic moment obtained by forcing all the
Mn moments lying in the (111) plane, and the calculated \( \sigma^2_{xy} \)
spectrum (not shown here) is nearly identical to that obtained
without fixing the Mn moments to lie in the (111) plane.

The dispersive parts of optical conductivity, i.e., \( \sigma^2_{xx} \) and
\( \sigma^2_{xy} \), can be obtained from the corresponding absorptive parts
by use of the Kramers-Kronig transformation. In Fig. 3 we
plot the \( \sigma^2_{xx} \) and \( \sigma^2_{xy} \) for the T1 and T2 spin configura-
tions, respectively. Figs. 3(a) and 3(b) show that, similar to the \( \sigma^2_{xx} \),
the \( \sigma^2_{xy} \) are almost the same for the T1 and T2 configurations.
The \( \sigma^2_{xy} \) for all the Mn\(_3\)X alloys have common characteris-
tics such as a broad valley around 1.0–1.5 eV, a negative to
positive crossing point at 2.5 eV, and a broad plateau above
3.0 eV. Figures 3(c) and 3(d) show that the \( \sigma^2_{xy} \) for the T1 and
T2 configurations have similar profiles but opposite in sign,
and gradually decay to small values in the high energy region,
which are similar to the behavior of \( \sigma^2_{xy} \).

C. Magneto-optical Kerr effect

After discussing the magnetic and optical properties of the
Mn\(_3\)X alloys, we now turn our attention to their magneto-
optical property. From the complex Kerr angle spectra pre-
sented in Fig. 4 one can find their key features as follows: (1)
The calculated Kerr rotation angles (\( \theta_K \)) and ellipticities (\( \epsilon_K \))
for the T1 and T2 states have opposite signs, inheriting from
the behaviors of the off-diagonal elements of optical conduc-
tivity, \( \sigma^1_{xy} \) [see Figs. 5(c) and 5(d)] and \( \sigma^2_{xy} \) [see Figs. 2(c)
and 5(d)].

FIG. 3. (Color online) (a) and (b) The calculated imaginary diagonal
(\( \sigma^2_{xx} \)) as well as (c) and (d) real off-diagonal (\( \sigma^1_{xy} \)) components
of optical conductivity tensor for the T1 and T2 spin configurations
of the Mn\(_3\)X alloys. Both \( \sigma^2_{xx} \) and \( \sigma^1_{xy} \) are broadened with a Lorentzian
of 0.1 eV to simulate the finite lifetime effect of electron.
and $(2d)$; (2) The sign reversals of $\theta_K$ and $\epsilon_K$ are frequent in the given energy range. When $\epsilon_K$ crosses the zero line, a peak turns up in the corresponding $\theta_K$ spectrum, and vice versa, which may be ascribed to the Kramers-Kronig relation; (3) The overall features and maximum values of $\theta_K$ and $\epsilon_K$ have a size sequence of Mn$_3$Pt $>$ Mn$_3$Ir $>$ Mn$_3$Rh, which suggests that the SOC strength of the $X$ atom contributes significantly to the enhancement of the MOKE spectra; (4) The largest Kerr rotation angle that appears in Mn$_3$Pt is $\sim 0.6^\circ$ at incident photon energy of 0.7 eV. This angle arising from a noncollinear antiferromagnet is remarkably large and comparable to those of transition metals, such as the bulk and multilayers of Fe and Co studied earlier.

Finally, we attempt to analyze the origin of the large MOKE in the Mn$_3$X alloys. Equation (3) indicates that a peak in a Kerr spectrum can stem from either a small $\sigma_{xx}$ in the denominator or a large $\sigma_{xy}$ in the numerator, which are called the “optical” and “magneto-optical” origins, respectively. From Figs. 3(c) and 4(a), one can observe that the positions of the peaks of $\sigma_{xy}^{T1}$ and $\theta_K$ for the T1 state overlap with each other, such as the first peak at 0.7 eV in Mn$_3$Pt and the twin peaks in the low energy range of 0~1.0 eV in Mn$_3$Ir, and the same can be seen for the T2 state by comparing Fig. 3(d) with Fig. 4(b). Furthermore, the positions of the peaks of $\epsilon_K$ and $\sigma_{xy}$ are close, as shown in Figs. 2(c) and 2(d) as well as Figs. 3(c) and 4(d), respectively. Thus, the magnitudes of the peaks of $\theta_K$ and $\epsilon_K$ are modulated by the $\sigma_{xx}$, as shown in Figs. 2(a) and 2(b) as well as Figs. 3(a) and 3(b), respectively. Since the Kerr rotation angle and also the ellipticity entangle in a complicated way with both the real and imaginary components of the optical conductivity tensor, there are no analytic forms for strictly separating the “optical” and “magneto-optical” origins for them. On the other hand, the nonzero $\sigma_{xy}$ is clearly the root cause for the emergence of the Kerr effect in this kind of noncollinear antiferromagnets, as already discussed in the preceding subsection. This is corroborated by our test calculations which show that both the $\sigma_{xy}$ and MOKE in these antiferromagnets would become zero without the SOC included. Therefore, it can be concluded that the large MOKE in the Mn$_3$X alloys has a “magneto-optical” origin rather than the “optical” origin.

IV. SUMMARY

In conclusion, using first-principles density functional calculations, we have investigated the possible magneto-optical Kerr effect in noncollinear antiferromagnets Mn$_3$Rh, Mn$_3$Ir and Mn$_3$Pt. We found that the Kerr rotation angle can be as large as $\sim 0.6^\circ$ in Mn$_3$Pt, which is comparable to that in elemental transition metal ferromagnets such as bcc Fe. We also discussed the differences in the magneto-optical responses for the T1 and T2 spin configurations. The surprisingly large magneto-optical Kerr effect in noncollinear antiferromagnets with nearly zero magnetization is attributed to the nontrivial off-diagonal components of optical conductivity, i.e., having the so-called “magneto-optical” origin. Our results demonstrate that one cannot assume a priori vanishing magneto-optical Kerr effect in antiferromagnets with zero net magnetization. The large Kerr rotation angle, plus other interesting physical properties of the Mn$_3$X alloys, such as prominent exchange-bias properties and high Néel temperature, would make these materials an exciting platform for exploring novel information-storage devices.

ACKNOWLEDGMENTS

Q.N. and G.Y.G. thank Hua Chen and Alan MacDonal for stimulating discussions. W.F. and G.Y.G acknowledge support from the Ministry of Science and Technology, the Academia Sinica and NCTS of Taiwan. W.F. and Y.Y. were supported in part by the MOST Project of China (Grant Nos. 2014CB920903 and 2011CBA00108), the NSF of China (Grant Nos. 11374033, 11225418, and 11174337), the Specialized Research Fund for the Doctoral Program of Higher Education of China (Grant Nos. 2012110110046 and 20131101120052), and the Basic Research Fund of Beijing Institute of Technology (Grant No. 20141842004). Q.N. was supported in part by DOE-DMSE (Grant No. DE-FG03-02ER45958) and the Welch Foundation (Grant No. F-1255). W.F. also acknowledges the use of the computational resources provided by the National Supercomputer Center in Tianjin (NSCC-TJ).

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