Time reversal and the spin angular momentum of transverse-electric and transverse-magnetic surface modes

Kyoung-Youm Kim,[1,*] Il-Min Lee,[2] Jungho Kim,[3] Jaehoon Jung,[4] and Byoungho Lee[2]

[1]*Department of Optical Engineering, Sejong University, Seoul 143-747, Korea*

[2]*National Creative Research Center for Active Plasmonics Application Systems, Inter-University Semiconductor Research Center and School of Electrical Engineering, Seoul National University, Seoul 151-744, Korea*

[3]*Department of Information Display, Kyung Hee University, Seoul 130-701, Korea*

[4]*Department of Electronics and Electrical Engineering, Dankook University, Gyeonggi-Do 140-714, Korea*

(Received 17 July 2012; revised manuscript received 12 November 2012; published 5 December 2012)

The polarization state of a surface mode is a hybrid between linear and elliptical states, meaning that although one of its electric- and magnetic-field components is linearly polarized, the other rotates on the transverse-longitudinal plane, effectively in its elliptical polarization state. We show that the rotation of the electric-field component can induce transverse spin angular momentum of the surface mode. The rotation of the magnetic-field component, however, cannot generate the spin. We argue that this results from the fact that the rotation direction of the magnetic field is invariant under time reversal.

DOI: 10.1103/PhysRevA.86.063805 PACS number(s): 42.25.Bs, 41.20.Jb, 78.20.Ci

Interfaces between two materials having opposite signs of constitutive parameters (permittivity ε and/or permeability μ) have attracted research interest because they can support a new class of waveguide mode, a *surface mode* [1]. One of many interesting features of these surface modes is that although their transverse wave numbers are imaginary and thus they seem to propagate only in their longitudinal directions, the optical power guided through them is transported in a vortexlike way [2]. That is, some of the optical power originally moving forward in one material layer is transported transversely (i.e., across the interface) and made to flow longitudinally again in the other layer, but in a backward direction. This vortexlike power flow makes it clear that the surface mode carries angular momentum (AM), as was recently shown independently by Kim and Lee [3] and by Bliokh and Nori [4].

Usually, optical waves can carry two different kinds of AM, the orbital AM and the spin AM [5]. They can be distinguished based on their interaction with matter, i.e., depending on whether they induce orbiting or spinning motions of the absorptive particles [6,7]. Light waves with helical phase fronts can carry the orbital AM [8] while their circular polarization states can induce the spin AM [9]. In this paper, we focus on the spin AM carried by surface modes. In particular, we pay attention to their polarization state, which is a hybrid between the linear and the elliptical states. That is, one of the electric- and magnetic-field components comprising the surface mode is linearly polarized but the other is in its elliptical polarization state, rotating on the transverse-longitudinal plane with passing time. We show that the rotation of the electric-field component can induce spin AM but that of the magnetic component cannot by considering whether their rotations are vector quantities like spin under time reversal.

Let us consider a material interface (with a normal vector $\hat{\mathbf{x}}$), which supports a surface mode along the $+z$ direction. In the case of TM polarization, the magnetic field of this mode can be written in a complex form as $\tilde{\mathbf{H}} = \hat{\mathbf{y}}\phi(x)\exp[i(\beta z - \omega t)]$, where ω is the angular frequency of light in vacuum and $\phi(x)$ $(= \phi_r + i\phi_i)$ denotes the transverse profile of the surface mode whose complex propagation constant is β $(= \beta_r + i\beta_i)$. In the real-field notation, we have

$$\begin{aligned}\mathbf{H} &= \mathrm{Re}\{\tilde{\mathbf{H}}\} \\ &= e^{-\beta_i z}\hat{\mathbf{y}}[\phi_r(x)\cos(\beta_r z - \omega t) - \phi_i(x)\sin(\beta_r z - \omega t)], \quad (1)\end{aligned}$$

where $\mathrm{Re}\{\cdot\}$ denotes the real part of the corresponding quantity. Throughout this paper, vector quantities with a tilde (such as $\tilde{\mathbf{H}}$) are in a complex form while those without a tilde are real vectors. The above mode field becomes that of the electric component ($\mathbf{E}$) if we assume TE light waves [see Eq. (8)]. For the calculation of the spin AM carried by the surface mode, we have to start with the linear momentum of light. Actually, there are two different forms of the linear momentum density of light in media: the Abraham and the Minkowski ones. The Abraham form is $\mathbf{g}^A = \mathbf{E} \times \mathbf{H}/c^2$ while the Minkowski form is $\mathbf{g}^M = \mathbf{D} \times \mathbf{B}$, where $\mathbf{D}$ and $\mathbf{B}$ are the electric displacement and the magnetic flux density, respectively. For the past hundred years, extensive efforts have gone into determining which of these is correct [10,11]. There remain strong theoretical arguments in favor of both momenta [12–14]. Experimental works have reported some evidence in favor of the Abraham momentum [15] but other evidence as well in support of the Minkowski momentum [16]. Only recently, it was pointed out [17] that both of them are correct, with the Abraham form being the kinetic momentum (the product of mass and velocity) and the Minkowski form the canonical momentum (Planck's constant divided by the de Broglie wavelength, associated with spatial translations). What appears in experiments that measure the effect of a phase shift (which is proportional to the spatial translation) or diffraction (which is connected to the wavelength) is the Minkowski momentum [17]. Hereafter, we will distinguish the quantities based on the Abraham momentum and those based on the Minkowski momentum using the superscripts A and M.

The Minkowski version of the total AM of light is given by $\mathbf{J}^M = \int \mathbf{r} \times [\mathbf{D} \times (\nabla \times \mathbf{A})]\, dV$ and can be written as a sum

*kykim@sejong.ac.kr

of its orbital (**L**) and spin (**S**) parts:

$$\mathbf{J}^M = \int \mathbf{r} \times \sum_j D_j \nabla A_{rot,j}\, dV + \int \mathbf{D} \times \mathbf{A}_{rot}\, dV = \mathbf{L}^M + \mathbf{S}^M. \quad (2)$$

The Abraham version is $\mathbf{J}^A = v_g^{-2} \int \mathbf{r} \times [\mathbf{E} \times \mathrm{Re}\{\mu^{-1}(\nabla \times \tilde{\mathbf{A}})\}]dV$, where $\mathbf{A} = \mathrm{Re}\{\tilde{\mathbf{A}}\}$ and v_g denotes the group velocity of light, defined as $(\partial\beta_r/\partial\omega)^{-1}$ in our case, resulting in

$$\mathbf{J}^A = \frac{1}{v_g^2} \int \mathbf{r} \times \sum_j E_j \mathrm{Re}\{\mu^{-1} \nabla \tilde{A}_{rot,j}\} dV + \frac{1}{v_g^2} \int \mathbf{E} \times \mathrm{Re}\{\mu^{-1} \tilde{\mathbf{A}}_{rot}\} dV = \mathbf{L}^A + \mathbf{S}^A. \quad (3)$$

Considering that surface modes are usually guided through dielectric and metamaterial layers, we can interpret these $\mathbf{L}^{M(A)}$ and $\mathbf{S}^{M(A)}$ as those carried by various types of *polaritons* excited in each material layer. That is, the spin AM of the surface mode is associated with the spinning of *quasiparticles* at the left and right sides of the interface and their resultant magnetic moment [18].

In deriving Eqs. (2) and (3), we have decomposed the vector potential **A** as $\mathbf{A} = \mathbf{A}_{rot} + \mathbf{A}_{irr}$, where $\mathbf{A}_{rot}$ and $\mathbf{A}_{irr}$ represent its rotational (whose divergence becomes zero) and irrotational (whose curl vanishes) parts, respectively. Although the vector potential itself is dependent on the gauge, $\mathbf{A}_{rot}$ is gauge invariant, causing the orbital and the spin AMs to be gauge invariant as well. If we choose the Coulomb gauge, $\mathbf{A}_{irr}$ reduces to zero and $\mathbf{A}_{rot}$ becomes equal to **A**, resulting in the following spin AM densities:

$$\mathbf{s}^M = \mathbf{D} \times \mathbf{A} = \mathrm{Re}\{\varepsilon\tilde{\mathbf{E}}\} \times \mathrm{Re}\{\tilde{\mathbf{A}}\} = \frac{\varepsilon_r}{\omega} \mathrm{Re}\{\tilde{\mathbf{E}}\} \times \mathrm{Im}\{\tilde{\mathbf{E}}\}, \quad (4)$$

$$\mathbf{s}^A = \frac{1}{v_g^2} \mathbf{E} \times \mathrm{Re}\{\mu^{-1}\tilde{\mathbf{A}}\} = \frac{1}{v_g^2|\mu|^2} \mathrm{Re}\{\tilde{\mathbf{E}}\} \times \mathrm{Re}\{\mu^*\tilde{\mathbf{A}}\} = \frac{\mu_r}{\omega v_g^2 |\mu|^2} \mathrm{Re}\{\tilde{\mathbf{E}}\} \times \mathrm{Im}\{\tilde{\mathbf{E}}\}, \quad (5)$$

where $\mathrm{Im}\{\cdot\}$ denotes the imaginary part of the corresponding quantity, and we have used $\tilde{\mathbf{E}} = i\omega\tilde{\mathbf{A}}$ or $\mathbf{E} = -\partial\mathbf{A}/\partial t$ since no sources are present in our case. We can easily calculate $\tilde{\mathbf{E}}$ of the TE and TM surface modes as follows:

$$\tilde{\mathbf{E}}_{TE} = e^{-\beta_i z} \hat{\mathbf{y}} \phi \exp(j\varphi_t), \quad (6)$$

$$\tilde{\mathbf{E}}_{TM} = \frac{e^{-\beta_i z}}{\omega\varepsilon} [\hat{\mathbf{x}}\beta\phi \exp(j\varphi_t) + \hat{\mathbf{z}} i\phi' \exp(j\varphi_t)], \quad (7)$$

where $\varphi_t = \beta_r z - \omega t$ and the prime indicates differentiation with respect to x. They give us

$$\mathbf{E}_{TE} = \mathrm{Re}\{\tilde{\mathbf{E}}_{TE}\} = e^{-\beta_i z} \hat{\mathbf{y}} [\phi_r \cos\varphi_t - \phi_i \sin\varphi_t], \quad (8)$$

$$\begin{aligned}\mathbf{E}_{TM} &= \mathrm{Re}\{\tilde{\mathbf{E}}_{TM}\} \\ &= \frac{e^{-\beta_i z}}{\omega|\varepsilon|^2} (-\hat{\mathbf{z}}[(\varepsilon_r\phi_i' - \varepsilon_i\phi_r')\cos\varphi_t \\ &\quad + (\varepsilon_r\phi_r' + \varepsilon_i\phi_i')\sin\varphi_t] + \hat{\mathbf{x}}\{[(\beta_r\varepsilon_r + \beta_i\varepsilon_i)\phi_r \\ &\quad + (\beta_r\varepsilon_i - \beta_i\varepsilon_r)\phi_i]\cos\varphi_t + [-(\beta_r\varepsilon_r + \beta_i\varepsilon_i)\phi_i \\ &\quad + (\beta_r\varepsilon_i - \beta_i\varepsilon_r)\phi_r]\sin\varphi_t\}). \end{aligned} \quad (9)$$

Using Eqs. (6) and (7), and the definitions of $\mathbf{s}^M$ and $\mathbf{s}^A$ in Eqs. (4) and (5), we can derive the following spin AM carried by TE and TM surface modes:

$$\mathbf{S}_{TE}^{M(A)} = 0, \quad (10)$$

$$\mathbf{S}_{TM}^{M(A)} = -\frac{e^{-2\beta_i z}}{\omega^3} \hat{\mathbf{y}} \int \Lambda^{M(A)} [\beta_r(\phi_r\phi_r' + \phi_i\phi_i') - \beta_i(\phi_r'\phi_i - \phi_r\phi_i')]\, dx, \quad (11)$$

where $\Lambda^M = \varepsilon_r/|\varepsilon|^2$ and $\Lambda^A = \mu_r/(v_g^2|\mu|^2|\varepsilon|^2)$. We have

$$\Lambda^A = \frac{\mu_r/\varepsilon_r}{v_g^2|\mu|^2} \cdot \Lambda^M, \quad (12)$$

where it can be easily seen that Λ^A and Λ^M have the same sign in dielectric [or positive-index medium (PIM)] and negative-index metamaterial (NIM) layers, but different signs in singly negative media such as ε-negative (ENG) and μ-negative (MNG) metamaterials. This entails that $\mathbf{s}_{TM}^A$ and $\mathbf{s}_{TM}^M$ are mutually parallel in PIM and NIM but become antiparallel in ENG and MNG materials. (Some more discussion will be given in the later part of this paper.)

We can observe two points in Eqs. (10) and (11). The first is that **S** is independent of the choice of the origin of the spatial coordinates. It is also independent of time, i.e., it is *constant* over time although the surface mode has a harmonic time dependency [$\exp(-i\omega t)$]. These characteristics confirm that it is an intrinsic quantity of the surface mode. The second point to note is that only the TM surface mode can carry the spin AM. The key principle behind the derivation of Eqs. (10) and (11) is that we can write the magnetic flux density **B** as a curl of the vector potential **A**. This, in turn, originates from the nonexistence of magnetic monopoles. Therefore, we can say that the fundamental difference in the characteristics of electric and magnetic fields, especially with regard to their monopoles, results in this breaking of duality. This interesting feature can be experimentally observed by using a Stern-Gerlach-like configuration with various scanning probe microscopy technologies such as the near-field scanning optical microscope (NSOM) or the photon scanning tunneling microscope (PSTM) [19]. Since the TM surface mode with a nonzero spin AM is associated with quasiparticles that have a nonzero magnetic moment, it will be deflected during propagation if we exert a magnetic-field gradient in a Stern-Gerlach way. By checking whether TE and TM surface modes undergo such a deflection or not, we can determine that they do or do not carry spin AM. In addition, by measuring the amount of the deflection or the transverse displacement, we can figure out the actual spin values carried by them. This measurement can also determine which form of the spin AM (Minkowski or Abraham) is involved. Actually, there is a very similar experiment which measured the deflection of a circularly polarized optical beam in a medium due to the acquired magnetic moment of the resultant polaritons [18]. However, there are two factors that can make our measurement more difficult. One is the short propagation length of the surface mode (due to the loss of metamaterials), which would result in a very small transverse displacement even if we use *slow* surface modes [20] as in [18]. The other, which is closely related to the above, is the change of the spin AM

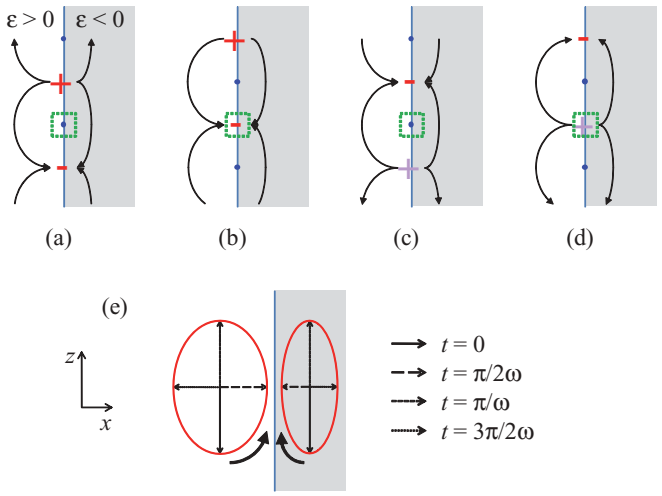

FIG. 1. (Color online) (a)–(d) Schematics showing the propagation of the electric-field component of a TM-polarized surface mode at $t = 0, \pi/2\omega, \pi/\omega$, and $3\pi/2\omega$, respectively. We assumed an interface between materials whose permittivities have opposite signs, which supports a forward mode. Peak positions of the induced (and accumulated) polarization charge density are denoted as $+$ and $-$. (e) Polarization states of the **E** field near the dotted area in (a)–(d).

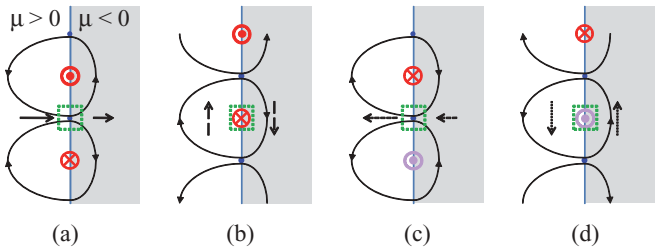

FIG. 2. (Color online) As Fig. 1 for the magnetic flux density (**B**) of a TE-polarized surface mode. We assumed similarly an interface between materials whose permeabilities have opposite signs, which also guides a forward mode. Peak positions of the induced (and accumulated) magnetization current density are denoted as $\odot$ and $\otimes$. Straight arrows near the dotted area indicate the local directions of the **B** field there.

during propagation [note that $\mathbf{S}_{TM}(z) = \mathbf{S}_{TM}(0)\exp(-2\beta_i z)$ in Eq. (11)] [3]. That is, the polaritons involved do not have a constant magnetic moment as in the usual Stern-Gerlach configuration. More research is now under way.

In what follows, we will investigate this polarization dependency of the spin AM further, focusing on the polarization states of TE and TM surface modes. In the case of a TM mode, although its magnetic-field component is always along the $\pm y$ directions, its electric component lies on the x-z plane, effectively in its elliptical polarization state. This is evident mathematically from Eq. (9), which gives us $\mathbf{E}_{TM} \propto \hat{\mathbf{x}}(\beta/\varepsilon)\cos\varphi_t - \hat{\mathbf{z}}(\kappa/\varepsilon)\sin\varphi_t$ in the lossless case where $\kappa = \pm\sqrt{\beta^2 - \varepsilon\mu\omega^2}$ and $\phi' = \kappa\phi$ [21]. However, we can draw a physical picture of this *rotation* of the **E** field by considering the induced polarization charges [22] given by $\rho_v = \varepsilon_0 \nabla \cdot \mathbf{E}$, where ε_0 is the permittivity in vacuum. In Fig. 1, we depict the temporal change of the **E** component of a TM surface mode, focusing on the induced polarization charges and their accumulation. As the peak positions of the accumulated polarization charge density (denoted as $+$ and $-$) move along the $+z$ direction as time goes on, the local directions of the **E** field change accordingly. For example, let us look into the area enclosed by (green) dotted lines. At $t = 0$, $\pi/2\omega$, π/ω, and $3\pi/2\omega$, the **E** field is along the $-z$, $+x$, $+z$, and $-x$ directions in the left layer, and along the $-z$, $-x$, $+z$, and $+x$ directions in the right layer, respectively [23]. These characteristics are summarized in Fig. 1(e), where we can see that the **E** field is in a kind of elliptical polarization state although its rotation directions are mutually opposite at the left and right sides of the interface.

On the contrary, in the TE mode case, the **B** field comes to be in the elliptical polarization state with the **E** field along the $\pm y$ directions. In Fig. 2, we described the propagation of the **B** field of a TE surface mode, similarly focusing on the induced magnetization currents [22] given by $\mathbf{J}_m = (1/\mu_0)\nabla \times \mathbf{B}$, where μ_0 is the permeability in vacuum. Both the **B** fields in the left and right layers are in the elliptical polarization state but their rotation directions are mutually opposite as in the case of Fig. 1(e).

We would like to argue that these rotations of the field components or their elliptical polarization states induce the spin AM carried by the surface modes. For this, however, we have to address two questions. The first is why only the rotation of the **E** field can generate the spin AM. The answer can be obtained by asking a different question: Is the spin even or odd under time reversal? That is, does the spin AM retain or reverse its original direction if time flows *backward*? Classical intuition lets us presume that the spin would change its direction, i.e., it is odd under time reversal. In Fig. 3(a), we plotted the polarization state of the **E**-field component of the TM surface mode analyzed in Fig. 1 under time reversal, taking Fig. 1(d) as a new initial field (at $t' = 0$). From the figure, we can see that the rotation directions of the **E** field are reversed. That is, the rotation direction of the **E** field is odd under time reversal. This is not the case, however, in the case of TE polarization. Let us look at Fig. 3(b) which is actually equal to Fig. 2(d). Under time reversal, the surface mode (i.e., the peak position of the induced magnetization current density) moves along the $-z$ direction, resulting in the field pattern shown in Fig. 3(c′). However, since the magnetization currents result from the movement of charges, their direction should be reversed as well, resulting in Fig. 3(c) instead of Fig. 3(c′). We note that Fig. 3(c) is actually equal to Fig. 2(a). That is, the same field pattern is produced irrespective of whether time flows forward or backward, which implies that the rotation direction of the **B** field is even or invariant under time reversal. The rotation of the **B** field is thus a kind of pseudovector under time reversal, making it impossible to be associated with vector quantities like the spin AM (which would require a sign flip). We argue that this is the reason why TE surface modes cannot carry the spin AM.

The second question concerns the opposite rotation directions of the **E** field at the left and right sides of the interface. If we neglect material losses for simplicity of discussion, the Minkowski and the Abraham forms of the spin AM density can be written as $\mathbf{s}^M = -\beta\phi\phi'/(\omega^3\varepsilon)\hat{\mathbf{y}}$ and $\mathbf{s}^A = -\beta\phi\phi'/(\omega^3 v_g^2 \mu\varepsilon^2)\hat{\mathbf{y}}$. The directions of $\mathbf{s}^M$ at the left and right sides of the interface are respectively proportional to

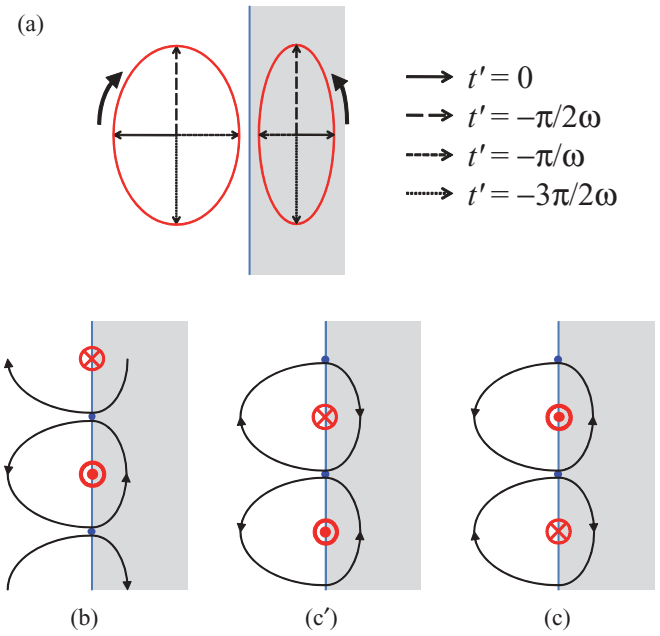

FIG. 3. (Color online) (a) As Fig. 1(e) under time reversal. We took Fig. 1(d) as a new initial field configuration ($t' = 0$). The rotation direction of the **E** field is reversed or is odd under time reversal. (b),(c) The case of TE polarization. We begin with (b) [which is equal to Fig. 2(d)]. Under time reversal, the surface mode propagates along the $-z$ direction, resulting in (c′). However, the direction of the magnetization current should be reversed as well, which produces the field pattern shown in (c). It is actually equal to Fig. 2(a). That is, the rotation direction of the **B** field is even under time reversal.

$-\beta/\varepsilon$ and β/ε while those of $\mathbf{s}^A$ are proportional to $-\beta/\mu$ and β/μ. In Fig. 4, we compared the directions of $\mathbf{s}^M$ and $\mathbf{s}^A$ at PIM-NIM, PIM-ENG, and MNG-ENG interfaces which guide forward TM modes. We can see that their directions are the same as the rotation direction of the **E** field in PIMs, but are opposite to that in NIMs [see Fig. 4(b)]. This difference can be easily understood by considering a circularly polarized light incident from PIM to impedance-matched NIM layers. From the continuity condition of the tangential **E** field and the fact that the phase velocity becomes antiparallel to the group velocity in a NIM, we can determine that the handedness of the circular polarization state of the incident light changes in a NIM (right-handed circular state to left-handed circular state and vice versa) while preserving its spin AM. This clearly indicates that the rotation direction of the **E** field in the NIM is opposite to the direction of the spin AM, and thus the different rotation directions of the **E** field at the left and right sides of the interface actually indicate the same direction of $\mathbf{s}^M$ or $\mathbf{s}^A$ as is shown in Fig. 4(b).

Unfortunately, however, we cannot think in this way in the cases of singly negative media since incident light waves cannot propagate through them. In Figs. 4(c) and 4(d), we can see that there are two directions of the spin AM density in

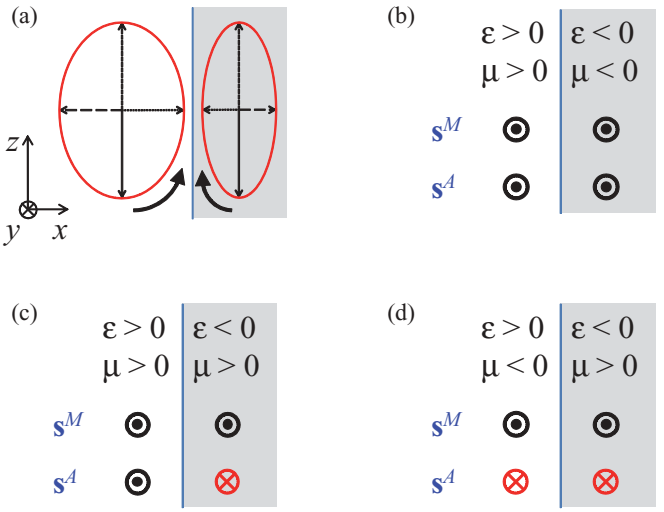

FIG. 4. (Color online) (a) Rotation directions of the **E** field at the left and right sides of the interface which supports a forward TM mode ($\beta > 0$). (b)–(d) Directions of the Minkowski and the Abraham forms of the spin AM density at various types of interface composed of (b) PIM-NIM, (c) PIM-ENG, and (d) MNG-ENG layers. One interface we neglected here is that between MNG and NIM layers. This is in order to avoid any confusion because it usually supports a backward mode, and in such a backward case, the directions of the spin AM (since $\beta < 0$) as well as those of the **E**-field rotation are all reversed [3].

singly negative layers, one being parallel with and the other antiparallel to the rotating direction of the **E** field. These two directions are respectively those of $\mathbf{s}^M$ and $\mathbf{s}^A$, and what is parallel with the rotation vector of the **E** field becomes $\mathbf{s}^A$ and $\mathbf{s}^M$ in ENG and MNG materials, respectively. As a result, we can observe that $\mathbf{S}^M$ and $\mathbf{S}^A$ carried by a TM surface mode at the MNG-ENG interface have opposite directions. In addition, $\mathbf{S}^A$ of the surface mode at a PIM-ENG interface can vanish while $\mathbf{S}^M$ does not. Further study from this perspective is necessary.

In conclusion, we showed that the surface mode which is supported not by the index difference but by the different signs of ε or μ across the interface can carry the spin AM. By investigating the polarization state of each field component comprising the surface mode, we showed that this spin AM originates from the rotation or the elliptical polarization state of the electric-field component of the TM surface mode. The TE surface mode, in which the magnetic component rotates instead, cannot carry the spin because the rotating direction remains invariant under time reversal, making its association with the spin (which requires a sign flip under time reversal) impossible.

This work was supported by the National Research Foundation and the Ministry of Education, Science and Technology of Korea through the Creative Research Initiative Program (Active Plasmonics Application Systems).

[1] B. E. Sernelius, *Surface Modes in Physics* (Wiley-VCH, Berlin, 2001).

[2] K.-Y. Kim, J. Kim, I.-M. Lee, and B. Lee, Phys. Rev. A **85**, 023840 (2012).

[3] K.-Y. Kim and B. Lee, in *Proceedings of the Second Japan-Korea Metamaterials Forum, June 2012*, edited by J. W. Wu and M. Hangyo (National Institute for Materials Science, Tsukuba, Japan, 2012), pp. 40–41.

[4] K. Y. Bliokh and F. Nori, Phys. Rev. A **85**, 061801(R) (2012).
[5] See S. Franke-Arnold, L. Allen, and M. Padgett, Laser Photonics Rev. **2**, 299 (2008); A. M. Yao and M. J. Padgett, Adv. Opt. Photon. **3**, 161 (2011) for recent reviews.
[6] A. T. O'Neil, I. MacVicar, L. Allen, and M. J. Padgett, Phys. Rev. Lett. **88**, 053601 (2002).
[7] V. Garcés-Chávez, D. McGloin, M. J. Padgett, W. Dultz, H. Schmitzer, and K. Dholakia, Phys. Rev. Lett. **91**, 093602 (2003).
[8] L. Allen, M. W. Beijersbergen, R. J. C. Spreeuw, and J. P. Woerdman, Phys. Rev. A **45**, 8185 (1992).
[9] R. A. Beth, Phys. Rev. **50**, 115 (1936).
[10] R. Peierls, *More Surprises in Theoretical Physics* (Princeton University Press, Princeton, NJ, 1991).
[11] P. W. Milonni and R. W. Boyd, Adv. Opt. Photon. **2**, 519 (2010).
[12] N. L. Balazs, Phys. Rev. **91**, 408 (1953).
[13] I. Brevik, Eur. J. Phys. **2**, 37 (1981).
[14] R. Loudon, S. M. Barnett, and C. Baxter, Phys. Rev. A **71**, 063802 (2005).
[15] W. She, J. Yu, and R. Feng, Phys. Rev. Lett. **101**, 243601 (2008).
[16] G. K. Campbell, A. E. Leanhardt, J. Mun, M. Boyd, E. W. Streed, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. **94**, 170403 (2005).
[17] S. M. Barnett, Phys. Rev. Lett. **104**, 070401 (2010).
[18] L. Karpa and M. Weitz, Nat. Phys. **2**, 332 (2006).
[19] A. Passian, A. L. Lereu, A. Wig, F. Meriaudeau, T. Thundat, and T. L. Ferrell, Phys. Rev. B **71**, 165418 (2005).
[20] K.-Y. Kim, C.-Y. Hwang, and B. Lee, Opt. Express **19**, 2286 (2011).
[21] Using the approach in [2], we can obtain qualitatively the same result (only the polarization ellipse is rotated with respect to the y axis) even if we consider the effect of material loss.
[22] K.-Y. Kim and J. Kim, Opt. Lett. **36**, 4065 (2011).
[23] This is because, as was pointed out in the caption of Fig. 1, we considered a forward mode. If a backward mode is assumed, the peak positions of the accumulated polarization charge density would move along the $-z$ direction, reversing the rotation directions of the **E** field in both layers.