Wireless at the Nanoscale: Toward Magnetically Tunable Beam Steering

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Abstract—Plasmonic wireless nanolinks hold great promise to overcome limitations from conventional metallic wires, namely, narrow bandwidths, Ohmic losses, dispersion, and cross-talking. However, current developments are limited to the wireless communication between two fixed points, i.e., a fixed transmitter with a fixed receiver, which in turn limits the energy efficiency and integration levels. In this work, we propose the use of magnetic fields for the active tuning of the radiation beam steering of plasmonic nanoantennas. The physical principle behind our concept is shown with a rigorous solution of the radiation pattern of a radiating dipole source in the presence of magnetic fields. The results indicate that the proposal of this work is feasible with magnetoplasmmonic nanoantennas, which can be fabricated using available ferromagnetic materials.

Index Terms—Beam steering, magnetooptic (MO) effects, magnetoplasmonic, nanoantennas.

I. INTRODUCTION

NANOSCALE analogs of conventional radio frequency (RF) antennas, commonly called nanoantennas, have become a subject of considerable attention during the past decade [1], [2], [3], [4], [5], [6], [7]. This idea was initially inspired by the possibility of using plasmonic dipoles as chip-scale optical/infrared wireless nanolinks [2], [3], with the potential to overcome the major drawbacks of conventional nanoelectronic interconnects (e.g., copper wires), namely, narrow bandwidths, Ohmic losses, dispersion, and cross-talking [8]. Furthermore, owing to the unique confinement properties of plasmonic structures, these nanoantennas hold promise for improved levels of integration density and miniaturization of on-chip photonic devices [9], [10]. However, there are two major limitations that need to be solved for this to become a reality. First, due to large optical wireless propagation losses, the transmitting (\(T\)) and receiving (\(R\)) nanoantennas are limited to be placed at distances in the order of the wavelength (\(\lambda\)). Second, in contrast to RF antennas, plasmonic wireless nanolinks are limited to the communication between two fixed \(T\) and \(R\), i.e., there are no mechanisms available to actively manipulate communication between a single \(T\) with multiple \(R\). Attempts to overcome the first limitation have mainly focused on geometric designs that optimize the values of directivity (\(D\)), which in turn optimize the ratio between the received (\(P_r\)) and transmitted (\(P_t\)) power (\(|P_r/P_t| \sim (D_t/D_r)^2\), where \(l\) represents the propagation length) [3]. Indeed, plasmonic horn-like [4], [5], Yagi–Uda-like [7], and plantenna [6] designs have demonstrated propagation distances of up to three orders of magnitude larger than the plasmon wavelength. In contrast, finding mechanisms to dynamically manipulate the beam steering of plasmonic nanoantennas remains a challenge.

The dynamic beam steering of plasmonic nanoantennas, using static magnetic fields, is demonstrated in this work. We refer to this latter mechanism as magnetically tunable beam steering. In contrast to 1-D and 2-D magnetooptic (MO) grating structures (with optical properties dictated by the corresponding diffraction orders), where the magnetic beam steering has been recently demonstrated [11], [12], we focused here on the active manipulation of the radiated field by a single magnetoplasmmonic nanoantenna. The physical principle behind our idea stems in the magnetic-field-induced optical anisotropies, widely known as MO effects. The latter depend on the direction of the applied static magnetic field (in relation to the direction of propagation of the wave) and on the polarization of the electromagnetic wave [13]. Significantly, MO effects have been used to tune the localized surface plasmon resonance (LSPR) features of isolated nanoparticles [14], [15], [16], as well as to control the light transmission, reflection, and absorption amplitudes by nanostructures made of noble metals [14], [17], ferromagnetic metals [15], [16], and hybrid systems [18].

II. THEORETICAL FRAMEWORK

A. Magnetooptics and Magnetic Materials

The propagation properties of electromagnetic waves are affected by the presence of a static magnetic field \(\mathbf{M}\), as it was first discovered by Faraday [19] and Kerr [20], [21],
respectively, in the mid-1800s. \( \mathbf{M} \) can be externally applied (e.g., using coils) due to the presence of the intrinsic magnetization of a ferromagnetic material. It was also observed that these effects are determined by the orientation of \( \mathbf{M} \) in relation to the polarization of the wave and its direction of propagation. For example, in the case of in-plane magnetized surface (i.e., \( \mathbf{M} \) along the plane of the surface), we can observe changes in the ellipticity and amplitude of the reflected/transmitted wave if \( \mathbf{M} \) is parallel or perpendicular to the plane of incidence (known as the longitudinal and transverse Kerr MO effects), respectively. On the contrary, for \( \mathbf{M} \) perpendicular to the surface, polarization rotation effects are observed for the reflected/transmitted waves (known as Faraday rotations) [22]. Indeed, these latter magnetization-induced polarization rotations [23] have been exploited for radar cross section reduction [24], [25] and magnetically tunable working polarization rotations [22]. Indeed, these latter magnetization-induced netization of a ferromagnetic material. It was also observed (e.g., using coils) due to the presence of the intrinsic magnetic field must remain applied netizable MO materials (e.g., graphene and InSb [30], [31], [32], where the external magnetic field must remain applied during device operation), ferromagnetic materials can retain their magnetized state after the external field is turned off. Moreover, the magnetization direction and sense can be tuned through the use of on-chip integrated coils [33] or using external electromagnetic coils [34]. The last mechanism, depicted in Fig. 1(a), is broadly used to manipulate magnetization of magnetoplasmonic nanostructures [34]. Different geometrical parameters of the nanoantenna design are shown in the upper and cross-sectional views in Fig. 1(b) and (c).

Although magnetic field effects can be considered in the relative permeability tensor \( \mu \) (when working in the RF domain), as in Pozar [24], [25] and Tan et al. [26], [27], in the optical regime it is conventionally used in the effective permittivity tensor (\( \varepsilon \)), taking \( \mu = 1 \). This last approach was rigorously demonstrated several decades ago [35]. Therefore, all the calculations in this work were carried out under the saturation magnetization condition, for which the permittivity values were used from the experimental data in Wang et al. [36]. It is worth remarking that the off-diagonal components (along \( xy \), \( xz \), or \( yz \)) of the permittivity tensor will depend on the orientation of \( \mathbf{M} \) (along \( z \), \( y \), or \( x \), respectively).

**B. Radiated Dipolar Fields in the Presence of a Static Magnetic Field**

Since the radiative properties of the nanostructure in Fig. 1 are dominated by dipolar features [28], we can start by idealizing it as an isolated point dipole \( \mathbf{p} \) (i.e., an infinitesimally small dipole) in the presence of \( \mathbf{M} \). As shall be demonstrated in Appendix, the magnetic beam steering is only achieved for \( \mathbf{M} \perp (\mathbf{p} \text{ and } \mathbf{k}_{\text{inc}}) \) (where \( \mathbf{k}_{\text{inc}} \) denotes the incident wave vector). Therefore, we use \( \mathbf{p} \), (i.e., \( \mathbf{p} \parallel \mathbf{x} \)) and \( \mathbf{M} \parallel \mathbf{y} \), with \( \mathbf{k}_{\text{inc}} \parallel \hat{z} \), as illustrated in Fig. 2(a). This orientation of \( \mathbf{M} \) induces an anisotropy along the \( xz \)-plane, described by the anisotropic dielectric tensor

\[
\hat{\varepsilon} = \begin{pmatrix} \varepsilon & 0 & im\varepsilon_{xz} \\ 0 & \varepsilon & 0 \\ -im\varepsilon_{xx} & 0 & \varepsilon \end{pmatrix}
\]  

(1)

where \( m = +1 \) (\( m = -1 \)) is used for \( \mathbf{M} \) pointing along the positive (negative) sense of the \( y \)-axis. \( \varepsilon_{xz} = \varepsilon_{zx} \) represents the contribution of the static magnetic field on the optical properties, which couples the \( x \)- and \( z \)-components of the optical electric field. In the case of ferromagnetic materials, as in this work, \( \varepsilon_{xz} \) is measured at the saturation magnetization for different wavelengths. In the case of nonferromagnetic materials such as, for instance, graphene and InSb, there are analytical expressions for \( \varepsilon_{xz} \propto H \) (\( H \) is the amplitude of the applied magnetic field) [30], [31], [32].

Using the point dipole polarizability in the radiative limit, \( \hat{\alpha} = (\hat{1} - i(k^3/6\pi)\hat{\alpha}_0)^{-1}\hat{\alpha}_0 \), with \( \hat{\alpha}_0 = 3V[\hat{\varepsilon} - \hat{1}][\hat{\varepsilon} + 2\hat{1}]^{-1} \) for the static (nonradiative) term, we obtain a polarizability tensor with the following structure:

\[
\hat{\alpha} = \begin{pmatrix} \alpha & 0 & m\alpha_{xz} \\ 0 & \alpha & 0 \\ -m\alpha_{zx} & 0 & \alpha \end{pmatrix}
\]  

(2)
where \( k \hat{z} \) and \( E \) are in vacuum, and fields, where \( k \) and \( E \) represent the radiated beam (from a dipolar source) can be actively tuned by the sign (along the \( y \)-axis) and magnitude of \( M \). More specifically, the sign and magnitude of \( M \) are in the \( m = \pm 1 \) and \( e \propto \varepsilon_{xx} \) terms in (8). This last result is numerically shown in Fig. 2(b), where the radiation pattern (at the \( xy \)-plane, i.e., \( \phi = 0^\circ \)) is presented for \( m = \pm 1 \). These numerical results are for a hypothetical point dipole with radius \( r_d = 90 \text{ nm} \), considering \( \lambda = 750 \text{ nm} \), \( e = -14 + i1.7 \), and \( \varepsilon_{xx} = -4.01 - i2.27 \). The magnetically tunable beam steering in Fig. 2(b) can be understood from the \( x \)-component of the scattered electric field profile (\( E_x^{\text{ sca}} \)), that is, the \( x \)-component from (3), which can be written as

\[
E_x^{\text{ sca}}(r, \theta, \phi) = \frac{k^3 E_0 \alpha e^{i kr}}{4 \pi (kr)^3} [-1 + kr(i + kr) + [-3 + kr(3i + kr)] \sin \theta \cos \phi (m\eta \cos \theta - \sin \theta \cos \phi)].
\]

The numerical results of (9) for \( m = \pm 1 \) are shown in Fig. 2(c) and (d) along the \( xy \)-plane, calculated at \( z = (\lambda/2\pi) \). The asymmetry of \( E_x^{\text{ sca}} \) for \( m = 1 \) and \( m = -1 \), around the axis \( x = 0 \), explains the change in the tilt direction of the radiation beam when the magnetic field is reversed.

### III. Results and Discussion

To demonstrate that the concept of magnetically tunable beam steering can be extended from idealized point dipoles to realistic magnetoplasmonic nanoantennas, let us now study the MO effects on the radiative properties of the nanoantenna in Fig. 1. Our idea is developed using a well-known nanos-structure, namely, a horn-like aperture plasmonic nanoantenna, which has been used for optical wireless communications [28] and plasmonic biosensing applications [39]. The system can be considered built by a noble metal (or ferromagnetic metal) film of thickness \( t \), with a horn-like nanoaperture, placed on a conventional \( \text{SiO}_2 \) glass substrate, as illustrated in Fig. 1. The nanoaperture has a rectangular cross section with a fixed length \( l_s \), whereas the width changes linearly from a length \( w_i \) at the top to a length \( w_i \) at the bottom of the structure [see Fig. 1(c)]. A rectangular cavity, of width \( w_i \) and depth \( d_e \), is used in the substrate (at the bottom of the nanoantenna) to optimize the resonance features of the system [28]. Following the analytical model described in (1)–(9), we consider the nanoantenna fed by an \( x \)-polarized incident plane wave, as illustrated in Fig. 1. The numerical results are obtained
to fabricate this type of structure, we used ion beam (FIB) milling is the most widely used technique [28].

In agreement with the current minimum size limitation of the ferromagnetic cobalt–silver alloy (Co\textsubscript{6}Ag) 

The resonant wavelength of the structure (maximum \(\sigma\)) in Fig. 3(a) and (b), respectively. Since focused ion beam (FIB) milling is the most widely used technique to fabricate this type of structure, we used \(w_1 = 10\) nm in agreement with the current minimum size limitation of the technique [28].

Focusing our attention on the structure with \(w_1 = 80\) nm and \(l_1 = 85\) nm, with a resonant wavelength \(\lambda = 700\) nm, we plot the normalized x-component of the scattered field (\(E_x^{sca}\)) along the \(xz\) - and \(yz\)-planes in Fig. 3(c) and (d), respectively. These calculations were performed for a demagnetized Co\textsubscript{6}Ag\textsubscript{94} film (\(M = 0\)), i.e., for \(m = 0\) in (1), where we used \(\epsilon\) and \(\epsilon_{xz}\) from the available experimental data [29]. In particular, we used \(\epsilon = -12.45 + i\,1.766\) and \(\epsilon_{xz} = -3.34721 - i\,1.75\) for \(\lambda = 700\) nm. The corresponding radiation patterns for \(m = 0\) and \(m = \pm 1\) are shown in Fig. 4(a) and (b). Although there is no consolidated model, the empirical equation \(G = (3 \times 10^3/\theta_{x}\theta_{y})\) (from RF counter-part) [42] is being actively used to estimate the maximum gain in optically driven nanoantennas [43]. In the latter expression, \(G\) represents the maximum gain, while \(\theta_{xz}\) and \(\theta_{yz}\) are the half-power beamwidth (HPBW) along the \(xz\) and \(yz\) elevation planes, respectively. Through this model, the demagnetized nanoantenna exhibits a linear gain \(G(m = 0) = 2.745\) for \(\theta = 0^\circ\). After magnetizing the system, the radiation pattern is slightly tilted from \(0^\circ\) to \(\theta = -3^\circ\) (\(\theta = +3^\circ\)) when \(m\) changes from 0 to +1 (−1). In contrast to Fig. 4(a) (for \(m = 0\), where \(G = 2.745\), we noted lower values of gain \((G(m = \pm 1) = 2.302)\) in Fig. 4(b), which are explained by larger optical losses (via the loss contribution of off-diagonal terms) and small magnetically tuned resonance changes (i.e., the resonance wavelength is changed slightly).

Indeed, magnetic tuning of plasmonic resonances is the most widely used principle for magnetoplasmonic biosensing applications [44], [45]. Importantly, as shown in Appendix, the magnetically tunable beam steering demonstrated in Fig. 4 can only be obtained for \(M\) placed perpendicular (transverse) to the plane formed by \(p_3\) and \(k_{nc}\) (illustrated in Fig. 2).

Using the geometric tuning of the resonant wavelength, we considered regions where the Co\textsubscript{6}Ag\textsubscript{94} film exhibits higher MO activity. In particular, we studied the system at \(\lambda = 725\) nm (\(\epsilon_{xz} = -3.681 - i\,1.174\)) and \(\lambda = 750\) nm (\(\epsilon_{xz} = -4.016 - i\,2.275\)), corresponding to \(w_1 = 160\) nm and \(w_2 = 220\) nm [shown in Fig. 3(a)], where the maximum changes in the beam steering were \(\theta_{xz} = \mp 4.6^\circ\) and \(\theta_{yz} = \mp 5.2^\circ\) (in relation to \(\theta = 0^\circ\)), respectively. The subindex \(\pm\) is used to indicate the sign of \(m\) in the calculations. The radiation patterns for \(w_1 = 220\) nm are shown in Fig. 5(a) and (b) for \(m = -1\) and \(m = +1\), respectively. It is important to note that in addition to enhanced beam steering, the gain is also enhanced \(G(m = \pm 1) = 2.95\) when compared with the case for \(w_1 = 80\) nm. The normalized \(E_x^{sca}\) near-fields associated with \(m = \mp 1\) are shown in Fig. 5(c) and (d).

We should remark that the numerical results for the magnetoplasmonic nanooantenna in Fig. 5 are in excellent qualitative agreement with the analytical results in Fig. 2 (for an isolated ideal dipole). Significantly, increasing \(\epsilon_{xz}\) would result in enhanced beam steering, as expected from (7) to (9) (where \(\epsilon_{xz}\) is implicit in \(\eta\)). In the case of ferromagnetic materials, the maximum value of \(\epsilon_{xz}\) is an intrinsic property of the material and, therefore, is limited by the corresponding saturation magnetization condition at different wavelengths. In contrast, for nonferromagnetic materials, \(\epsilon_{xz}\) is proportional to the amplitude of the applied magnetic field [30], [31], [32].
Interestingly, graphene and InSb materials exhibit strong MO effects [30], [31], [32] at terahertz (THz) frequencies, where they have also been used to design and develop THz antennas [46], [47], [48], [49]. Hence, we can hope that the tunable effects [30], [31], [32] at terahertz (THz) frequencies, where they have also been used to design and develop THz antennas [46], [47], [48], [49]. Hence, we can hope that the tunable effects. Hence, we can hope that the tunable antennas demonstrated here can be implemented not only in the visible and near-infrared ranges but also at the THz frequency range.

IV. CONCLUSION

Summarizing, we have theoretically shown that magnetic fields can be used to dynamically tune the direction of the beam radiated by plasmonic nanoantennas (which we called, in general, magnetoplasmonic nanoantennas). It was analytically and numerically demonstrated that the applied magnetic field should be placed perpendicular to the plane formed by the dipole ($\textbf{p}_d$) and the incident wavevector ($\textbf{k}_{\text{inc}}$). Our concept provides a new paradigm for active beam steering of wireless signals at nano- and micro-scale levels. From the applications point of view, our idea enables the communication of a single $T_c$ with multiple $R_c$, therefore allowing for improved energy efficiency and higher integration levels (through a reduced number of on-chip/interchip elements). Besides, the physical principle in this work lies on the optical radiation properties of electric dipolar sources in the presence of a magnetic field. Hence, the proposed scheme is not limited to a specific plasmonic nanoantenna design, i.e., it can be applied to any plasmonic antenna from the visible to the THz range.

APPENDIX

All the calculations in this work were performed considering that the radiating electric dipole is placed along the $x$-axis, as described in the main text. Here, we are only showing the absence of magnetically tunable beam steering for the other two possible configurations of $\textbf{M}$, which are $\textbf{M} \parallel \hat{x}$ and $\textbf{M} \parallel \hat{z}$.

1) Let us start discussing the case $\textbf{M} \parallel \hat{x}$, where the permittivity and polarizability tensors read

$$
\begin{align*}
\hat{\varepsilon} &= \begin{pmatrix}
\varepsilon & 0 & 0 \\
0 & \varepsilon & m\varepsilon_{yz} \\
0 & -m\varepsilon_{yz} & \varepsilon
\end{pmatrix}, \\
\hat{\alpha} &= \begin{pmatrix}
\alpha & 0 & 0 \\
0 & \alpha & m\alpha_{yz} \\
0 & -m\alpha_{yz} & \alpha
\end{pmatrix}.
\end{align*}
$$

Therefore, we have the dipolar moment $\textbf{p} = \varepsilon E_0 |\alpha| \hat{x}$ and radiated density power

$$
S_{k,\text{rad}}(\textbf{r}) = \frac{1}{2Z_0} \left( k^3 |E_0| |\alpha| \right)^2 \frac{f(\theta, \phi)}{(2k\lambda)^2}
$$

with $f(\theta, \phi) = 3 + \cos 2\theta - 2\sin^2 \theta \cos 2\phi$.

From (12), we can directly observe that neither $|\textbf{M}|$ nor $m$ will alter the corresponding radiation pattern of dipolar sources. The latter was corroborated with full electromagnetic simulations for the magnetoplasmonic nanoantenna, as shown in Fig. 6. These numerical results indicate a small diminishing in the amplitude of the radiation pattern, when compared with 2.745 for $m = 0$, which can be understood in terms of small changes in the effective refractive index of the magnetoplasmonic medium.

2) For the case of $\textbf{M} \parallel \hat{z}$, the permittivity and polarizability tensors are written as

$$
\begin{align*}
\hat{\varepsilon} &= \begin{pmatrix}
\varepsilon & 0 & i\varepsilon_{xy} \\
0 & i\varepsilon_{xy} & \varepsilon \\
0 & 0 & \varepsilon
\end{pmatrix}, \\
\hat{\alpha} &= \begin{pmatrix}
\alpha & 0 & 0 \\
0 & \alpha & 0 \\
0 & 0 & \alpha
\end{pmatrix}.
\end{align*}
$$

Fig. 5. Gain radiation patterns for the system with $w_x = 220$ nm and $l_x = 85$ nm for (a) $m = -1$ and (b) $m = +1$. Red dotted lines are used to indicate the radiated beam direction for each $m$. (c) and (d) Corresponding scattered near-fields for $m = -1$ and $m = +1$, respectively.

Fig. 6. Numerical results for the (a) and (b) near- and (c) and (d) far-fields associated to $\textbf{M} \parallel \hat{x}$ for $m = \pm 1$, where the subindex $x$ is used to emphasize the different orientations of $\textbf{M}$ in relation to the main text. Calculations were made for the structure in Fig. 4 of the main text, for $\lambda = 700$ nm, only changing the orientation of $\textbf{M}$.
changing the orientation of \( M \) in relation to the main text. Calculations were made for the structure in Fig. 4 of the main text, for \( \lambda = 700 \text{ nm} \), only changing the orientation of \( M \).

Fig. 7. Numerical results for the (a) and (b) near- and (c) and (d) far-fields associated to \( M \parallel \hat{z} \) for \( m_z = \pm 1 \), where the subindex \( z \) is used to emphasize the different orientations of \( M \) in relation to the main text. Calculations were made for the structure in Fig. 4 of the main text, for \( \lambda = 700 \text{ nm} \), only changing the orientation of \( M \).

and

\[
\hat{\alpha} = \begin{pmatrix} \alpha & m \alpha_{xy} & 0 \\ -m \alpha_{xy} & \alpha & 0 \\ 0 & 0 & \alpha \end{pmatrix}. \tag{14}
\]

Therefore, the radial component of the Poynting vector is given by

\[
S_{k,\text{rad}}(r) = \frac{1}{2\varepsilon_0} \left( \frac{k^3 |E_0| |\alpha|}{4\pi} \right)^2 \frac{f(\theta, \phi)}{(2kr)^2} \tag{15}
\]

where

\[
f(\theta, \phi) = (|\eta|^2 + 1)(3 + \cos 2\phi) + 2(|\eta|^2 - 1)\sin^2 \theta \cos 2\phi + 4m\text{Re}[\eta] \sin^2 \theta \sin 2\phi \tag{16}
\]

with \( \eta = \alpha_{xy}/\alpha \).

Equations (15)–(16) indicate that neither \(|M|\) nor \(m\) can be used to manipulate the beam steering of the radiated field. However, the amplitude of \( S_{k,\text{rad}} \) depends on both \(|M|\) and \(m\). This last effect can be observed in Fig. 7, where full electromagnetic simulations were performed for the magnetoplasmonic nanoantenna in Fig. 4 of the main text, but changing the orientation of \( M \) to \( \hat{z} \). To understand this last effect, we should remember that \( M \parallel k_{\text{inc}} \) produces a Faraday rotation effect on the radiated beam, as can be noted from Fig. 8, which also slightly changes the resonant properties of the nanoantenna.

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