Lateral Chirality-sorting Optical Spin Forces in Evanescent Fields

Amaury Hayat*
School of Engineering and Applied Sciences, Harvard University, Cambridge, Massachusetts 02138, USA and École Polytechnique, Palaiseau 91120, France

J. P. Balthasar Müller† and Federico Capasso
School of Engineering and Applied Sciences, Harvard University, Cambridge, Massachusetts 02138, USA

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The unusual transverse component of the spin angular momentum of evanescent waves gives rise to lateral forces on chiral particles, which have the surprising property of acting in a direction in which there is neither a field gradient nor wave propagation. The direction of these forces is opposite for particles with opposite helicities, such that they may be useful for optically-induced enantiomer separation with a single beam, and the reliance on an evanescent field makes them a natural choice for sorting within an integrated optical circuit. The magnitude of these forces substantially exceeds those of the recently predicted sideways optical forces acting on non-chiral objects in evanescent fields and on chiral objects in propagating fields near a surface, such that they may more readily offer an experimental confirmation of lateral optical forces.

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I. INTRODUCTION

Chirality (from Greek χειρ (cheir), 'hand') is a type of asymmetry where a geometric object can not be made to coincide with its mirror image through any proper spatial rotation [1]. From elementary particles, chiral molecules and crystals to human hands, chiral objects abound in nature, and the ongoing study of chirality-dependent effects has attracted considerable interest in chemistry, biology and physics ever since the discovery of optical activity two centuries ago. A chiral object and its mirror image are in some cases difficult to tell apart as they share all properties other than their handedness, yet they differ substantially in their behavior in a chiral environment such as plants or the human body [2, 3]. In the creation of chiral compounds usually both helicities are initially present, such that their analysis and separation represents a significant manufacturing challenge in research and industry. This for example includes the pharmaceutical industry, where enanto-pure drugs account for a major share of today’s world market [4], and chirality-dependent pharmacology and pharmacokinetics require the close monitoring of chirality [5]. While the sorting of molecules and materials by chirality normally has to be addressed through the introduction of a specific chiral resolving agent [6], the manifestation of chirality in the electromagnetic response of materials has raised the possibility of passive sorting using optical forces [7, 8]. Chirality-dependent optical forces emerge in particular because electromagnetic waves can themselves be chiral, which has been shown to have a range of interesting consequences for light-matter interactions [10–12].

In Ref. [9] Wang et. al. recently predicted an electromagnetic plane wave to exert a lateral optical force on a chiral particle above a reflective surface, which emerges as the particle interacts with the reflection of its scattered field. This highly unusual force acts in a direction in which there is neither wave propagation nor an intensity gradient, and deflects particles with opposite helicities towards opposite sides. Apart from its fundamental interest, such a force may in theory be useful for scalable all-optical enantiomer sorting with a single, homogenous beam. In Ref. [13], Bliokh et. al. almost simultaneously predicted another lateral force that is exerted on non-chiral particles in an evanescent wave. This force is a consequence of the linear momentum of the wave that is associated with its spin angular momentum, which was first described by F. J. Belinfante in the context of quantum field theory, and which vanishes for a propagating plane wave [14–15].

The spin angular momentum itself, however, gives rise to a linear momentum transfer when light interacts with chiral materials. Consequently, engineering the local spin angular momentum density of optical fields can be used to taylor chirality-dependent optical forces. The aim of this work is to show that this effect results in particularly strong lateral forces as chiral particles interact with the extraordinary spin angular momentum of an evanescent wave. The magnitude of these forces substantially exceeds those of the previously predicted lateral forces, which may more readily enable an experimental test of the surprising phenomenon of transverse optical forces. The force also has the effect of laterally deflecting chiral particles according to their chiral polarizability (Fig. 1), such that it has the potential to be used for a convenient optical scheme for sorting and measuring material chirality without an optical lattice.

In the following section of this paper we provide a brief general discussion of the optical forces exerted by an electromagnetic wave on a small particle, and subsequently treat the specific case of an evanescent wave in the third
section. Finally, we provide estimates for the magnitude of the force that may be achieved for different types of small particles that are commonly discussed in the literature: molecules, spherical nanoparticles and nanohelices.

II. OPTICAL FORCES ON SMALL PARTICLES

The optical forces exerted on an object is rigorously calculated by first finding the distribution of electromagnetic fields and then integrating the Maxwell stress tensor over a surface enclosing the object. The scattering problem associated with finding the field distributions is tremendously simplified in the dipole approximation, which holds in the so-called Rayleigh limit that applies to particles much smaller than the wavelength. The dipole solution is sufficiently accurate for most molecular scattering problems and is furthermore the basis for efficient numerical methods that are used to find solutions for complex larger objects \(^{10}\). We consider the optical forces on a small particle in a source-free, lossless, non-dispersive and isotropic medium with relative electric permittivity \(\varepsilon\) and relative magnetic permeability \(\mu\). In the dipole approximation, the time-averaged total optical force \(\mathbf{F}\) exerted on the particle by a monochromatic electromagnetic wave is given by \(^{17}\):

\[
\mathbf{F} = \frac{1}{\pi} \text{Re} \left\{ \mathbf{d} \left( \nabla \otimes \mathbf{E}^* \right) + \mathbf{m} \left( \nabla \otimes \mathbf{H}^* \right) \right\} - \frac{k^2}{3} \text{Re} \left\{ \frac{1}{\mu} \left( \mathbf{d} \times \mathbf{m}^* \right) \right\},
\]

(1)

where \(\mathbf{E}\) and \(\mathbf{H}\) are the electric and magnetic field vectors of the incident electromagnetic wave at the location of the particle; \(\mathbf{d}\) and \(\mathbf{m}\) are the electric and magnetic dipole moments of the particle; \(k = n \omega / c\) is the wavenumber of the electromagnetic wave; \(\omega\) is the frequency; \(c\) is the speed of light in vacuum; and \(n = \sqrt{\varepsilon\mu}\) the refractive index of the medium. The fields are written in complex phasor notation throughout this paper, where the factor \(\exp(-i\omega t)\) giving the time dependence is implied, and the superscript * denotes the complex conjugate. Vector quantities are indicated by bold letters, and all expressions are given in Gaussian units unless stated otherwise. \(\hat{x}, \hat{y}, \text{ and } \hat{z}\) are unit vectors along the corresponding coordinate axes. The symbol \(\otimes\) denotes the dyadic product, so that the terms of the form \(\mathbf{W} (\nabla \otimes \mathbf{V})\) in eqn. \(^{1}\) have elements \(\left[ \mathbf{W} (\nabla \otimes \mathbf{V}) \right]_i = \sum_j W_j \partial_i V_j\) for \(i, j \in \{x, y, z\}\) \(^{17}\) and can be written as \(\mathbf{W} (\nabla \otimes \mathbf{V}) = (\mathbf{W} \cdot \nabla) \mathbf{V} + \mathbf{W} \times (\nabla \times \mathbf{V})\) in terms of more commonly used vector operators.

Equation \(^{1}\) shows that the optical force has a component \(\mathbf{F}_0\) corresponding to the force exerted by the incident field on the particle’s electric and magnetic dipole moments, and a component \(\mathbf{F}_{\text{int}}\) that results from a direct interaction of the two dipole moments. In case of the evanescent field, it is the \(\mathbf{F}_{\text{int}}\) term that contains the lateral electromagnetic spin force.

Chirality manifests itself in the electromagnetic response of a material through a cross-coupling of the induced electric and magnetic polarizations, such that an electric field also gives rise to a magnetic polarization and vice versa. Indeed, the optical response of chiral materials may be modeled intuitively by a simple coupled oscillator model \(^{18}\). The electric and magnetic dipole moments \(\mathbf{d}\) and \(\mathbf{m}\) induced by fields incident on a chiral particle with isotropic polarizabilities and no permanent dipole moment are then given by \(^{9}\):

\[
\mathbf{d} = \alpha_e \mathbf{E} + i \chi \mathbf{H}, \quad \mathbf{m} = \alpha_m \mathbf{H} - i \chi \mathbf{E},
\]

(2)

where \(\alpha_e\) and \(\alpha_m\) are the dynamic electric and magnetic polarizabilities. The chiral polarizability of the particle \(\chi\) captures the chiral nature of the dipole, such that setting \(\chi = 0\) recovers the case of an achiral dipole.

It is worth mentioning that, while the chiral polarizability is an inherently dynamic quantity, the dynamic dipole polarizabilities \(\alpha_e\) and \(\alpha_m\) differ from the more frequently listed corresponding static polarizabilities \(\alpha_e^{(0)}\) and \(\alpha_m^{(0)}\) by a radiation correction first derived by Draine as \(\alpha_e = \left[ 1 - \left( \frac{i 2 k^3}{3 \varepsilon} \right) \alpha_e^{(0)} \right]^{-1} \alpha_e^{(0)}\) and \(\alpha_m = \left[ 1 - \left( \frac{i 2 k^3}{3 \mu} \right) \alpha_m^{(0)} \right]^{-1} \alpha_m^{(0)}\) \(^{19}\). While the static polarizability is sufficiently accurate in the particular case of the small gold sphere considered in ref. \(^{13}\), omission of the radiation correction can in general lead to a substantial underestimation of the optical forces \(^{20}\).

We can combine eqns. \(^{1}\) and \(^{2}\) to rewrite the optical force on a chiral particle in terms of field quantities,
Here we marked the terms corresponding to the familiar gradient force and radiation pressure, both of which are modified in the presence of chirality. The final term of eqn. 3 shows that the spin momentum density of the field gives rise to a linear momentum transfer on chiral particles, which occurs in opposite directions for opposite signs of \( \chi \), i.e. opposite helicities. The prefactor \( g = (g/2\omega) \) arises from the usage of Gaussian units, \( u_e = (ge/4)|E|^2 \) and \( u_m = (gm/4)|H|^2 \) are the energy densities of the electric and the magnetic field, and \( h = (g/2\omega) \) is the time-averaged optical helicity density. The time-averaged Poynting momentum density \( p = (g/2\omega)Re \{ E \times H^* \} \) can be decomposed as \( p = p^s + p^o \) into the components related to the orbital angular momentum \( l \) and the spin angular momentum \( s \) according to \( l = r \times p^o \) and \( s = r \times p^s \), where \( p^s \) is the Belinfante spin momentum density. These momenta and angular momenta can then be individually further separated into their electric and magnetic components, in particular \( p^o = p_{e}^o + p_{m}^o \) and \( s = s_{e} + s_{m} \) (the supplementary material contains a full list of the definitions of the field quantities). Involving a fluid mechanics analogy, the term \( \nabla \times p \) may be interpreted as the vorticity of the photon flow. \( p'' = (g/2\omega)Im \{ E \times H^* \} \) is the imaginary Poynting momentum. For a monochromatic wave, the momentum and spin densities are:

\[
p^s = -\frac{g}{8\omega} \nabla \times \left[ (i\mu)^{-1} E \times E^* + (i\epsilon)^{-1} H \times H^* \right]
\]

\[
p^o = -\frac{g}{4\omega} Im \left\{ \mu^{-1} E (\nabla \otimes E^*) + \epsilon^{-1} H (\nabla \otimes H^*) \right\}
\]

\[
s = -\frac{g}{4\omega} \left\{ (i\mu)^{-1} E \times E^* + (i\epsilon)^{-1} H \times H^* \right\}
\]

where the magnetic and electric contributions correspond to the terms proportional to the electric and magnetic fields. Evaluating eqns. (5) - (7) for a plane wave \( E = \sqrt{\mu} E_0 \exp(ikz) \) propagating in the \( z \)-direction, the only non-vanishing quantities leading to optical forces are:

\[
p/2 = p^o/2 = p_{e}^o = p_{m}^o = \frac{g}{4\omega} ik \hat{z}
\]

\[
s/2 = s_{e} = s_{m} = \frac{g}{4\omega} I \sigma \hat{z}
\]

which gives:

\[
F_0 = \frac{1}{g} \left( \epsilon^{-1} Re \{ \alpha_e \} \nabla u_e + \mu^{-1} Re \{ \alpha_m \} \nabla u_m \right) - \frac{\omega}{g} Re \{ \chi \} \nabla h + \frac{2\omega}{g} \mu Im \{ \alpha_e \} \mathbf{p}_e^o + \epsilon Im \{ \alpha_m \} \mathbf{p}_m^o - \frac{\epsilon}{g} Im \{ \chi \} (\nabla \times \mathbf{p} - 2k^2 \mathbf{s})
\]

\[
F_{\text{int}} = -\frac{2\omega k^3}{g} \left( Re \{ \alpha_e \alpha_m^* \} \mathbf{p} - Im \{ \alpha_e \alpha_m^* \} \mathbf{p}'' + |\chi|^2 \mathbf{p} \right) = \frac{2\omega^2 k^4}{g} \left( Re \{ \chi m \} \mathbf{s}_m + \mu Re \{ \chi e \} \mathbf{s}_e \right).
\]

where \( I = n^{-1} |E \times H^*| = |E_0|^2 \) is the intensity. The factor \( \sigma = -2Im \left\{ (E \cdot \hat{x})(E^* \cdot \hat{y}) \right\}/(E \cdot \hat{x})^2 + |E \cdot \hat{y}|^2 \) is a measure for the degree of circular polarization (ellipticity) of the wave in the (x,y)-plane, such that \( \sigma = +1 \) for right circular polarization and \( \sigma = -1 \) for left circular polarization. In the calculations that follow, we will furthermore use the parameters \( \tau = (E \cdot \ hat{x})^2 - (E \cdot \ hat{y})^2 / (E \cdot \ hat{x})^2 + |E \cdot \ hat{y}|^2 \) and \( \xi = 2Re \left\{ (E \cdot \hat{x})(E^* \cdot \hat{y}) \right\} / (E \cdot \hat{x})^2 + |E \cdot \hat{y}|^2 \) as measures for the linear polarization in the (x,y)-plane, such that \( \tau = \{+1,-1\} \) correspond to linear polarization along \( x,y \) and \( \xi = \{+1,-1\} \) correspond to linear polarization at \(+45^\circ,-45^\circ\) with respect to the \( x \)-axis. The forces on a dipole in a plane wave are purely in the direction of propagation:

\[
F_0 = \left[ \left( Im \{ \mu \alpha_e \} + Im \{ \epsilon \alpha_m \} \right)/2 + Im \{ \chi \} \sigma \right] \hat{k} \hat{z}
\]

\[
F_{\text{int}} = -\left[ \left( Re \{ \chi m \} + \mu Re \{ \chi e \} \right)/\sigma + Re \{ \alpha_e \alpha_m^* \} + |\chi|^2 \right] \hat{k}^4/3 \hat{z}
\]

The helicity-dependent change of the radiation pressure due to the chiral polarizability \( \chi \) has previously been used for sorting highly chiral liquid crystal droplets in optical lattices.

### III. OPTICAL FORCES IN AN EVANESCENT FIELD

We now may consider an evanescent wave created by the total internal reflection of light at an interface in the (x,y)-plane, which has the form \( E = \sqrt{\mu} E_0 \exp(-\kappa x) \exp(i k z) \). If the evanescent field is propagating in a medium with index \( n_2 = \sqrt{\epsilon} \) and was created through the total internal reflection of a beam in a medium with index \( n_1 \) incident on the interface with angle \( \theta \) in the (x,z)-plane (Fig. 1), then the wave vector components are given by \( k_x = n_1/n_2 \sin \theta \) \( k \), \( k_z = \sqrt{k_x^2 - k_0^2} \) and \( \kappa = \sqrt{k_x^2 - k_z^2} = i k_z \), where \( k = n_2 \omega/c \). The expressions for all quantities relevant for the calculation of the forces with eqns. 3 and 4 are listed in the supplementary information to save space. Notably, evanescent fields have longitudinally polarized field
components [20], which give rise to transverse spin angular momentum. In case of a p-polarized wave \((\tau = 1)\) it is the electric field that is elliptically polarized in the \((x,z)\)-plane, and in case of a s-polarized wave \((\tau = -1)\) it is the magnetic field. Correspondingly, the out of plane electric and magnetic spin-components are given by:

\[
(\hat{y} \cdot \mathbf{s}_e) = (1 + \tau) \frac{g}{4\omega} \frac{\kappa}{k_z} I_e \\
(\hat{y} \cdot \mathbf{s}_m) = (1 - \tau) \frac{g}{4\omega} \frac{\kappa}{k_z} I_c.
\] (12) (13)

It follows for the forces on a chiral dipole in an evanescent field that:

\[
\begin{align*}
\mathbf{F}_0 &= \kappa I_e \Re\{\chi\} n_2 \sigma - \frac{1}{2} \left( \Re\{\mu \alpha_e\} (1 + \tau) \frac{\kappa^2}{k_z} + \Re\{\epsilon \alpha_m\} (1 - \tau) \frac{\kappa^2}{k_z} \right) \hat{x} \\
&\quad - \frac{I_e}{2} k_z \Im\{\mu \alpha_e\} (1 + \tau) \frac{\kappa^2}{k_z} + \Im\{\epsilon \alpha_m\} (1 - \tau) \frac{\kappa^2}{k_z} + 2\Im\{\chi\} \sigma n_2 \hat{z}
\end{align*}
\] (14)

\[
\begin{align*}
\mathbf{F}_{\text{int}} &= -\frac{I_e}{3} \frac{k^4}{k_z} \frac{\kappa}{k_z} \Re\{\mu \alpha_e \alpha_m^*\} \sigma + \Im\{\alpha_e^* \alpha_m\} \xi + \sigma |\chi|^2 + \frac{\epsilon}{n_2} (1 - \tau) \Re\{\chi \alpha_m^*\} + \frac{\mu}{n_2} (1 + \tau) \Re\{\chi \alpha_e^*\} \hat{y} \\
&\quad - \frac{I_e}{3} \frac{k^4}{k_z} \Re\{\alpha_e^* \alpha_m\} |\chi|^2 + \frac{\sigma}{n_2} (\epsilon \Re\{\chi \alpha_m^*\} + \mu \Re\{\chi \alpha_e^*\}) \hat{z} \\
&\quad - \frac{I_e}{3} \frac{k^4}{k_z} \frac{\kappa k}{k_z^2} (\frac{\xi}{n_2} (\epsilon \Re\{\mu \chi \alpha_m^*\} - \Re\{\epsilon \chi \alpha_m^*\}) + \Im\{\alpha_e^* \alpha_m\} \tau) \hat{x}
\end{align*}
\] (15)

with the intensity of the evanescent field \(I_e = \frac{|E_0|^2}{1 + (\sigma/\kappa)^2} \exp(-2\kappa x)\). The lateral forces experienced by the dipole are given by the first line of eqn. [15] with all chirality-dependent terms proportional to \(\chi\). Unfortunately the chirality-dependent lateral forces caused by \(\nabla \times \mathbf{p}\) and \(\mathbf{s}\) in for the \(\mathbf{F}_0\) force exactly cancel each other. The \(y\)-component of the Poynting momentum density \(\mathbf{p}\) arises entirely from the Belinfante spin momentum density \(\mathbf{p}^s\) for polarizations with non-vanishing \(\sigma\), which also gives rise to a helicity-independent lateral force term on chiral particles. An analogous \(y\)-term arises as part of the imaginary Poynting momentum density \(\mathbf{p}^i\) for polarizations with non-vanishing \(\xi\). The intensity and polarization of the evanescent wave depend on the complex transmission coefficients for the totally internally reflected beam. To illustrate the characteristic magnitudes of the force, we use eqn. [15] to calculate the chirality-dependent lateral optical forces arising in several different systems.

IV. LATERAL FORCES ON CHIRAL PARTICLES

A. Helicene Molecules

The optical manipulation of objects smaller than a few 100 nm in liquid suspension is in general very challenging due to thermal agitation. However, the optical sorting of individual atoms or molecules by chirality may nevertheless be possible in low-pressure environments or by tailoring the fields in modern metamaterials. A molecule of hexahelicene has a real chiral polarizability of \(\chi \approx 4 \times 10^{-21} \text{A}^3\) at a wavelength of 589.3 nm, which lies well outside of the molecule’s absorption band.[22] At the same wavelength, the static electric polarizability of hexahelicene is \(\alpha_e = 10.4 \text{A}^3\) with the magnetic polarizability \(|\alpha_m| < 10^{-5}|\alpha_e|\). We consider a beam with a given intensity that is totally internally reflected at the interface of heavy flint glass \((n_1 = 1.74)\) and air \((n_2 = 1)\) at an angle of \(\theta = 36.4^\circ\), which is close to the critical angle and gives the strongest force given the angle-dependent intensity and wave-vector of the evanescent field. Situated 60 nm above the interface, the hexahelicene molecule experiences a lateral force of \(F_y \approx 2.04 \cdot 10^{-19} \text{pN/}(\text{mW/\mu m})\) if \(\sigma = 1\) (circularly polarized in the (x,y)-plane) and \(F_y \approx 4.04 \cdot 10^{-19} \text{pN/}(\text{mW/\mu m})\) if \(\tau = 1\) (p-polarized). Stronger optical forces are possible within the absorption band, where the chiral polarizability is complex and highly dependent on the wavelength. Coronene is a close relative of hexahelicene which is achiral but has the same size and electric polarizability as helicene. The lateral force on this achiral molecule due to the Belinfante momentum density \(\mathbf{p}^s\) as predicted by Blok et. al. [12] is \(F_y \approx 1.7 \cdot 10^{-21} \text{pN}\) for \(\sigma = 1\) and vanishes for \(\sigma = 0\) (linearly polarized in (x,y) plane).
B. Spherical Nanoparticles

The polarizability of a chiral particle may be calculated using bi-isotropic constitutive relations for the material of the particle, where material chirality is parametrized with a chirality parameter \( K \in [-1,1] \) [3]. The derivation and the somewhat bulky expressions for the polarizabilities of a small chiral sphere are given in the supplementary information. An example for spherical particles with highly chiral electromagnetic response are cholesteric liquid crystal droplets [8]. Assuming the refractive index of 5CB (4-Cyano-4′-pentylbiphenyl) of \( n_2 = 1.597 \) and \( K = 1 \) for a 30 nm sphere in water with index \( n_2 = 1.33 \), letting an incident beam with 589 nm wavelength totally internally reflect at \( \theta = 52^\circ \) in heavy flint glass (\( n_1 = 1.74 \)) causes a lateral force of \( F_y \simeq 2.6 \cdot 10^{-5} \text{pN}/(\mu \text{W}/\mu \text{m}^2) \) if the sphere is floating 60 nm above the flint glass surface and \( \tau = 1 \).

The lateral force on a chiral sphere above a reflective surface in a propagating field as predicted by Wang et. al. [9] corresponds to \( F_y = 10^{-8} \text{pN}/(\mu \text{W}/\mu \text{m}^2) \) for a sphere with 30nm radius in vacuum with a dielectric constant of \( \epsilon = 2 \) and chirality parameter \( K = 1 \) when it is located 60 nm above a metal surface and illuminated with a plane wave that has \( \tau = 1 \) and a wavelength of 600 nm. In comparison, the same sphere would experience chirality-sorting lateral force of \( F_y = 4.6 \cdot 10^{-5} \text{pN}/(\mu \text{W}/\mu \text{m}^2) \) at the same height above the surface in an evanescent field that is created by total internal reflection at \( \theta = 36.4^\circ \) in heavy flint glass (\( n_1 = 1.74 \)). Fig 2a shows a the magnitude of the lateral forces exerted on a chiral and an achiral spherical Rayleigh particle as predicted by Bliokh et. al. and Wang et. al. in comparison to the lateral force predicted here.

C. Gold Helix

Increasing interest in chiral optical materials has recently driven the development of artificial nanostructures with extreme chiral responses [23,20], including artificially chiral effective media [18]. Metallic nanohelices in particular enable the analytical estimation of their very strong chiral polarizability and have recently been fabricated at size scales small enough to reach plasmonic resonances in the visible range [24]. A perfectly conducting 50nm helix with 25nm diameter and one loop has a chiral polarizability of \( \chi = 2.72 \cdot 10^7 \text{Å}^3 \) and a static electric polarizability of \( \alpha_e = 2.56 \cdot 10^8 \text{Å}^3 \) [27]. Situated 30nm above the surface in an evanescent field with wavelength 589nm and \( \tau = 1 \), a lateral force of \( F_y = -4.25 \cdot 10^{-3} \text{pN}/(\mu \text{W}/\mu \text{m}^2) \) results if the beam is totally internally reflected at a flint glass / air interface at a \( \theta = 36.4^\circ \) degree angle.

V. CONCLUSION

We predict lateral forces on materials with chiral optical response in evanescent fields, which push particles with opposite helicities in opposite directions. The forces result from the interaction of the evanescent field’s transverse optical spin angular momentum density of the wave with the chiral electromagnetic response of the particle and are particularly strong in comparison to previously predicted lateral optical forces. Evanescent fields arise whenever light is laterally confined, so that the effect described in this paper represents a natural choice for the optical sorting of material chirality in an integrated system. The effect can be created by a single beam, which significantly aids alignment and avoids standing wave patterns that limit the separation of enantiomers to the width of interference fringes. However, the use of multiple beams may enable the cancellation of the longi-
tudinal force component or, given the possibility of generating transverse spin, even the generation of spin-optical lateral forces in free space beams [31]. For the sake of length and clarity, we limited the discussion to the simplest case of an evanescent field created by the total internal reflection of light at a single interface, but just as in conventional optical tweezer there is significant potential for enhancing the magnitude of the force by engineering the optical fields. Natural choices for this are the intense and inherently evanescent fields of plasmonic excitations, or optical waveguides designed to have regions with high field enhancement, such as slot waveguides [29]. We furthermore limited the discussion to the Rayleigh limit, where compact closed form expressions exist. By confining the discussion to particles in the dipole approximation, we were able to clearly show the relationship of the emerging optical forces to the field quantities of the electromagnetic field, though larger particles (such as in the Mie regime) are bound to experience much stronger forces.

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