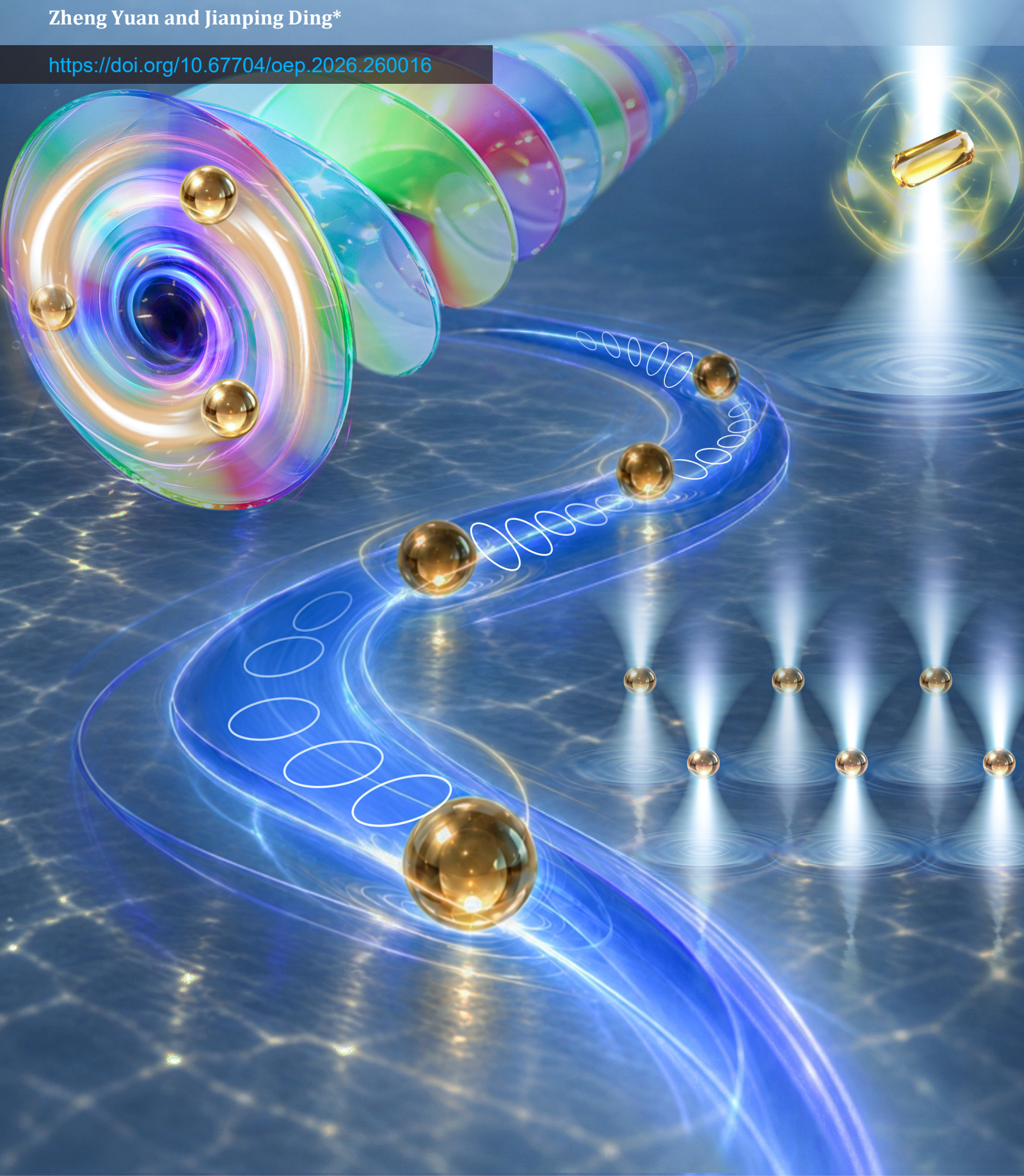


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Spatially structured light for optical micromanipulation

Zheng Yuan^{1,2,3} and Jianping Ding^{1,3,4*}

Abstract: Optical micromanipulation enables the noncontact trapping and motion control of micro- and nanoscale objects through linear and angular momentum exchange between light and matter. Conventional Gaussian optical tweezers provide limited control over spatial field structure, restricting their capabilities in parallel multiparticle manipulation, complex-trajectory transport, and multidimensional motion control. By tailoring amplitude, phase, and polarization, spatially structured light can reshape intensity gradients, local energy flow, and momentum and angular momentum distributions, enabling the on-demand design of optical forces, torques, and trapping potentials. This review introduces the orbital–spin decomposition of optical momentum and discusses how its real and imaginary components relate to the electromagnetic responses of particles. Representative scalar and vector structured-light modes are then surveyed, together with advances in field generation from two-dimensional complex-amplitude modulation to three-dimensional scalar and fully vectorial field synthesis. Recent developments in multiparticle trapping, complex-trajectory transport, dark-trap manipulation, particle rotation and orientation, three-dimensional control, and light-driven micro- and nanorobots are reviewed, with emphasis on emerging mechanisms associated with polarization gradients, spin-related momentum, and imaginary momentum. Future directions include spatiotemporal and ultrafast structured light, dynamically reconfigurable devices, diverse structured waves, and the manipulation of quantum and biological objects. By clarifying the connections among spatial field structure, electromagnetic momentum transfer, and particle motion, this review provides a basis for developing high-dimensional, dynamic, and integrated optical micromanipulation systems.

Keywords: spatially structured light; optical micromanipulation; optical tweezers; optical momentum; angular momentum; vector optical fields

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1 Introduction

Optical fields carry energy, linear momentum, and angular momentum. The scattering, refraction, and absorption of light by a particle alter the fluxes of electromagnetic linear and angular momentum. The resulting momentum transfer manifests as optical forces and torques that drive the translational and rotational motion of the particle^{1–3}. On this basis, optical tweezers can trap, translate, and rotate micro- and nanoscale objects without mechanical contact^{2–4}. Since Ashkin first used radiation pressure to propel particles and subsequently demonstrated a single-beam gradient-force optical trap, optical tweezers have evolved into an impor-

tant platform for investigating light–matter interactions, microscopic mechanical responses, and biophysical processes^{2–6}. They have been widely applied to measurements of minute forces and torques^{7–10}, single-molecule force spectroscopy^{11–13}, microrheology and soft-matter physics^{14–19}, nanostructure manipulation^{20–21}, studies of cells and biological systems^{22–24}, nonequilibrium statistical physics^{25–27}, and spectroscopic analysis and sorting^{28–30}.

Conventional single-beam optical tweezers typically employ a Gaussian beam tightly focused by a high-numerical-aperture objective^{3,4}. The steep intensity gradient near the focus draws particles with a refractive index higher than that of the surrounding medium toward the high-intensity

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region. A stable three-dimensional optical trap is formed when the axial gradient force counterbalances the scattering force directed along the beam propagation direction^{3,5}. Owing to its simple optical configuration, high stability, and well-established calibration procedures, this arrangement remains the fundamental architecture of optical micromanipulation^{6,8}. Nevertheless, the trapping potential generated by a single Gaussian focus has a relatively fixed geometry and is primarily suited to single-point trapping and focal scanning. For tasks such as the parallel manipulation of multiple particles^{31,32}, transport along curved trajectories^{33–34}, particle rotation and orientation^{9–10,35}, and control using random optical potentials^{36–37}, merely translating or multiplexing Gaussian foci makes it difficult to simultaneously achieve a high degree of control, a broad dynamic range, and efficient use of optical power.

Spatially structured optical fields provide a powerful route toward expanding the capabilities of optical micromanipulation^{38–40}. As illustrated in Fig. 1, tailoring the spatial distributions of amplitude, phase, and polarization enables control over the optical-intensity gradient, local energy flow, and

linear- and angular-momentum densities within the interaction region, thereby reshaping optical forces, optical torques, and trapping potentials^{38–42}. According to the spatial distribution of their polarization states, structured optical fields are commonly classified as scalar or vector fields. Scalar fields have an approximately uniform polarization state across the transverse plane, with their spatial structure residing primarily in the amplitude and phase. In vector fields, the orientation or ellipticity of the polarization ellipse varies with position, making polarization itself a spatial degree of freedom available for optical-field design^{40–41}.

Scalar structured fields include holographic optical fields for the parallel trapping of multiple particles^{31–32}, vortex beams for ring-shaped trapping and orbital motion^{43–45}, Bessel beams for long-range axial manipulation^{46–47}, Airy beams for transport along curved trajectories^{48,49}, and speckle fields for studies of random optical potentials and diffusive dynamics^{36,37,50}. Uniform circularly or elliptically polarized light can also induce the rotation of birefringent and anisotropic particles through the transfer of spin angular momentum^{9,10,35}. Vector structured light further intro-

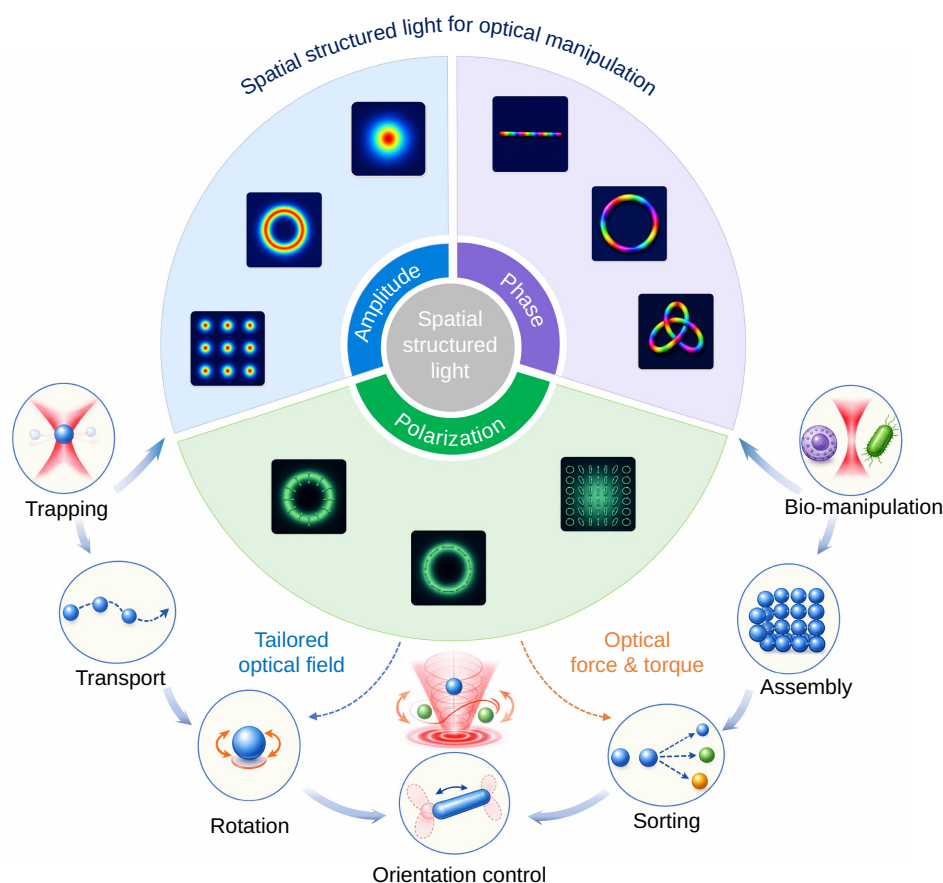


Fig. 1 | Overall framework of spatially structured optical fields for optical micromanipulation. The spatial manipulation of amplitude, phase, and polarization jointly determines the intensity gradient, local energy flow, and distributions of linear and angular momentum in the interaction region. Through absorption, scattering, and multipole responses of the particles, these field quantities are converted into applications such as trapping, transport, rotation, orientation, sorting, and assembly.

duces spatially nonuniform polarization, producing richer distributions of local spin density, transverse momentum, and their associated curls^{51–56}, while providing additional opportunities for the selective manipulation of polarization-sensitive and chiral objects^{57–59}. Importantly, amplitude, phase, and polarization do not exhibit simple one-to-one correspondences with the gradient force, orbital angular momentum, and spin angular momentum, respectively. Tight focusing, optical scattering, multipolar coupling, particle geometry and material response, and the surrounding boundaries can all modify the field distribution and momentum exchange within the actual interaction region^{5,56,60–61}.

This review focuses on the physical connections between spatially structured optical fields and optical micromanipulation, together with the technological evolution of this field. Section 2 introduces the fundamental descriptions of linear and angular momentum and discusses how spatial structures in amplitude, phase, and polarization are converted into mechanical action through the electromagnetic response of a particle. Section 3 reviews representative scalar and vector modes, as well as methods for generating and controlling two- and three-dimensional structured optical fields. Section 4 summarizes the applications of structured optical fields to trapping, transport, rotation, manipulation of low-refractive-index particles, multidimensional control, and light-driven micro- and nanorobots. Finally, Section 5 discusses the principal challenges facing the field and outlines promising directions for future development.

2 Fundamental Principles

2.1 Optical Momentum Density and Its Orbital–Spin Decomposition

Consider a monochromatic electromagnetic field of angular frequency ω , with the time dependence $e^{-i\omega t}$ and complex field amplitudes $\mathbf{E}(\mathbf{r})$ and $\mathbf{H}(\mathbf{r})$. In a homogeneous, isotropic, and nondispersive medium, the complex momentum density can be defined as

$$\Pi = \frac{1}{2c^2} \mathbf{E}^* \times \mathbf{H}, \quad (1)$$

where $c = 1/\sqrt{\epsilon\mu}$ is the speed of light in the medium, and ϵ and μ are the permittivity and permeability, respectively. The real part, $\mathbf{P} = \text{Re}(\Pi)$, corresponds to the time-averaged Minkowski momentum density^{42,55}. This momentum density can be decomposed into orbital and spin components, denoted by $\mathbf{P}_o$ and $\mathbf{P}_s$, respectively^{42,62}:

$$\begin{aligned} \mathbf{P} &= \mathbf{P}_o + \mathbf{P}_s \\ &= \frac{1}{4\omega} \text{Im} [\epsilon \mathbf{E}^* \cdot (\nabla) \mathbf{E} + \mu \mathbf{H}^* \cdot (\nabla) \mathbf{H}] + \frac{1}{2} \nabla \\ &\quad \times \frac{1}{4\omega} \text{Im} [\epsilon \mathbf{E}^* \times \mathbf{E} + \mu \mathbf{H}^* \times \mathbf{H}]. \end{aligned} \quad (2)$$

The orbital momentum density is given by

$$\begin{aligned} \mathbf{P}_o &= \mathbf{P}_o^e + \mathbf{P}_o^m \\ &= \frac{\epsilon}{4\omega} \text{Im} [\mathbf{E}^* \cdot (\nabla) \mathbf{E}] + \frac{\mu}{4\omega} \text{Im} [\mathbf{H}^* \cdot (\nabla) \mathbf{H}], \end{aligned} \quad (3)$$

whereas the spin momentum density takes the form

$$\begin{aligned} \mathbf{P}_s &= \mathbf{P}_s^e + \mathbf{P}_s^m \\ &= \frac{\epsilon}{8\omega} \nabla \times \text{Im} [\mathbf{E}^* \times \mathbf{E}] + \frac{\mu}{8\omega} \nabla \times \text{Im} [\mathbf{H}^* \times \mathbf{H}], \end{aligned} \quad (4)$$

Here, $\mathbf{P}_o^e$ and $\mathbf{P}_o^m$ denote the electric and magnetic contributions to the orbital momentum density, while $\mathbf{P}_s^e$ and $\mathbf{P}_s^m$ represent the corresponding contributions to the spin momentum density. The notation used in Eqs. (2) and (3) is defined by $\mathbf{C}^* \cdot (\nabla) \mathbf{C} = C_x^* \nabla C_x + C_y^* \nabla C_y + C_z^* \nabla C_z$.

This decomposition treats the electric and magnetic fields on an equal footing, reflecting what Berry termed “electric–magnetic democracy”⁴². In many practical analyses, Maxwell’s equations can be used to express the same momentum decomposition entirely in terms of the electric field:

$$\mathbf{P} = \mathbf{P}_o + \mathbf{P}_s = \frac{\epsilon}{2\omega} \text{Im} [\mathbf{E}^* \cdot (\nabla) \mathbf{E}] + \frac{1}{2} \nabla \times \frac{\epsilon}{2\omega} \text{Im} [\mathbf{E}^* \times \mathbf{E}], \quad (5)$$

Equations (3)–(5) reveal two types of momentum sources with distinct physical properties in structured optical fields. The orbital momentum density $\mathbf{P}_o$ describes the spatial energy redistribution within an optical field in the course of its evolution. For example, in a paraxial beam, the longitudinal variation of the optical energy density along the propagation direction (z -axis) is entirely determined by the divergence of the transverse orbital momentum density⁶².

The spin momentum density $\mathbf{P}_s$, by contrast, originates from spatial variations in the local polarization state or, equivalently, from the inhomogeneity of the spin angular momentum density. It describes the locally circulating energy flow. A useful intuitive picture is to regard a circularly polarized field as an assembly of microscopic “vortex units,” each carrying a local circulation determined by the polarization helicity. In a spatially uniform field, the circulations of adjacent vortex cells cancel each other, so no macroscopic spin flow exists. However, when the intensity or polarization state varies within the field, this cancellation becomes incomplete, leaving a residual spin flow near the beam boundary or in regions where the polarization changes abruptly, as shown in Fig. 2.

2.2 Contributions of Amplitude, Phase, and Polarization to Optical Momentum and Angular Momentum

To clarify how the amplitude, phase, and polarization structure of a light field contribute to its momentum density, the complex electric-field amplitude can be written as

$$\mathbf{E}(\mathbf{r}) = u(\mathbf{r}) e^{i\psi(\mathbf{r})} \hat{\mathbf{e}}(\mathbf{r}), \quad (6)$$

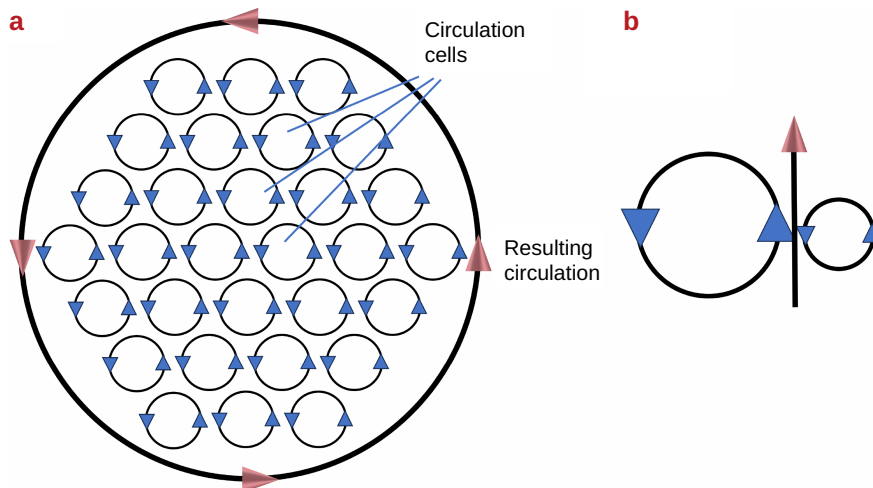


Fig. 2 | Schematic of spin-flow generation in a circularly polarized beam. (a) Uniform beam with a sharp boundary. (b) In an inhomogeneous beam, the imbalance of local cell circulations (indicated by different cell sizes) gives rise to a macroscopic spin flow (vertical arrow).

where $u(\mathbf{r})$ is a real-valued amplitude, $\psi(\mathbf{r})$ is the spatial phase, and $\hat{\mathbf{e}}(\mathbf{r})$ is a normalized complex polarization vector satisfying $\hat{\mathbf{e}}^* \cdot \hat{\mathbf{e}} = 1$. The corresponding intensity is denoted by $I(\mathbf{r}) = u^2(\mathbf{r})$. With this representation, the momentum density can be separated into three contributions with distinct physical origins⁶³:

$$\mathbf{p} = \underbrace{\frac{\epsilon}{2\omega} I(\mathbf{r}) \nabla \psi}_{\mathbf{p}_0} - i \underbrace{\frac{\epsilon}{2\omega} I(\mathbf{r}) \nabla \epsilon_j(\mathbf{r})}_{\mathbf{p}_s} + \frac{1}{2} \nabla \times \mathbf{s}(\mathbf{r}) \quad (7)$$

The first term arises from the spatial phase gradient $\nabla \psi$ and constitutes the conventional phase-gradient contribution to the orbital momentum density. It also underlies the familiar phase-gradient force exerted by structured optical fields³³. The second term arises from the spatial variation of the polarization vector $\hat{\mathbf{e}}$. It represents the contribution of the phase difference between the two components of the polarization state to the orbital momentum density. Together, these two terms form $\mathbf{P}_O$. The third term is the curl of the spin angular momentum density and constitutes the spin momentum density $\mathbf{P}_S$.

The spin angular momentum density is defined as⁶³

$$\mathbf{s}(\mathbf{r}) = \frac{\epsilon}{2\omega} I(\mathbf{r}) \boldsymbol{\sigma}(\mathbf{r}), \quad \boldsymbol{\sigma}(\mathbf{r}) = i \hat{\mathbf{e}}(\mathbf{r}) \times \hat{\mathbf{e}}^*(\mathbf{r}), \quad (8)$$

where $\boldsymbol{\sigma}(\mathbf{r})$ characterizes the local polarization helicity⁶³. For a paraxial beam propagating along the z direction, $\boldsymbol{\sigma} \cdot \hat{\mathbf{z}}$ corresponds to the normalized Stokes parameter S_3/S_0 . The right- and left-circular polarization correspond to $\boldsymbol{\sigma} \cdot \hat{\mathbf{z}} = +1$ and -1 , respectively.

The momentum density can further be used to define the angular momentum density, $\mathbf{J} = \mathbf{r} \times \mathbf{p}$, whose longitudinal component is $J_z = (\mathbf{r} \times \mathbf{p}) \cdot \hat{\mathbf{z}}$. The Jones vector of an arbitrary polarization state can then be parameterized as follows⁶⁴:

$$\hat{\mathbf{e}}(\mathbf{r}) = \begin{pmatrix} \cos \chi \cos \phi - i \sin \chi \sin \phi \\ \cos \chi \sin \phi + i \sin \chi \cos \phi \end{pmatrix}, \quad (9)$$

where $\chi(\mathbf{r})$ is the ellipticity angle and $\phi(\mathbf{r})$ is the orientation angle of the polarization ellipse. In cylindrical coordinates (ρ, φ, z) , the longitudinal angular momentum density can be expressed as⁶⁵

$$J_z = \frac{\epsilon}{4\omega} [2I(\mathbf{r}) \partial_\varphi \psi(\mathbf{r}) - 2I(\mathbf{r}) \sin 2\chi(\mathbf{r}) \partial_\varphi \phi(\mathbf{r}) - \rho \partial_\rho (I(\mathbf{r}) \sin 2\chi(\mathbf{r}))], \quad (10)$$

The three terms in Eq. (10) have different physical origins^{63,65–66}. The first arises from the azimuthal gradient of the spatial phase and represents the conventional orbital angular momentum contribution. For example, a vortex phase of the form $\psi = m\varphi$ directly produces an orbital angular momentum contribution proportional to the topological charge m .

The latter two terms are associated with the vectorial structure of the optical field. The second term originates from the azimuthal variation of the polarization-ellipse orientation in a vector beam and represents a spin-orbit conversion contribution. The third term arises from the radial variation of $I(\mathbf{r}) \sin 2\chi(\mathbf{r})$, thereby describing the contribution associated with radial gradients in both intensity and polarization ellipticity. For a locally linearly polarized vector field, $\chi = 0$, and the latter two terms vanish. In this case, the angular momentum density is governed primarily by the spatial phase structure. In an elliptically polarized vector field, however, the polarization-dependent terms may provide substantial and non-negligible contributions.

In addition to its real part, the imaginary part of the complex momentum density also carries physical information. It can be defined as⁶⁷

$$\mathbf{P}_I = \text{Im}(\Pi) = \frac{1}{2c^2} \text{Im}(\mathbf{E}^* \times \mathbf{H}). \quad (11)$$

For a locally linearly polarized field with an arbitrary polarization direction, the polarization vector $\hat{\mathbf{e}}$ may be

chosen to be real. If it further satisfies $\nabla \times \hat{\varepsilon} = 0$, Eq. (6), together with Maxwell's equation $\nabla \times \mathbf{E} = i\omega\mu\mathbf{H}$ gives⁶⁸

$$\mathbf{P}_I = -\frac{\varepsilon}{2\omega} \text{Re} [E^* (\nabla E - \hat{\varepsilon} (\hat{\varepsilon} \cdot \nabla E))] . \quad (12)$$

Using $I = EE^*$, Eq. (12) can be rewritten as

$$\mathbf{P}_I = -\frac{\varepsilon}{4\omega} [\nabla I - \hat{\varepsilon} (\hat{\varepsilon} \cdot \nabla I)] . \quad (13)$$

Equation (13) shows that the imaginary momentum density is independent of the spatial phase ψ . Instead, it is governed by the intensity gradient and its orientation relative to the local polarization direction. Consequently, even a locally linearly polarized field can possess a nonzero imaginary momentum density when the intensity distribution is spatially nonuniform.

The amplitude, phase, and polarization of a structured optical field jointly determine its local momentum and angular momentum distributions. In optical micromanipulation, however, the actual motion of a particle is ultimately governed by how these field quantities are converted into mechanical action through the particle's electromagnetic response and scattering. The following section therefore examines the fundamental optical-force mechanisms from the perspectives of the Maxwell stress tensor and the dipole approximation.

2.3 Optical-Force Mechanisms in Structured Light Fields

A rigorous evaluation of optical forces generally begins with the Maxwell stress tensor. By integrating the stress tensor over a closed surface Σ enclosing the particle, the total optical force can be written as^{4,69–70}

$$\mathbf{F} = \oint_{\Sigma} \langle \mathbf{T} \rangle \cdot \mathbf{n} dA , \quad (14)$$

where $\mathbf{n}$ is the outward unit normal to the surface and dA is the surface element. The time-averaged Maxwell stress tensor is given by

$$\langle \mathbf{T} \rangle = \frac{1}{2} \text{Re} \left[\varepsilon \mathbf{E} \mathbf{E}^* + \mu \mathbf{H} \mathbf{H}^* - \frac{1}{2} (\varepsilon |\mathbf{E}|^2 + \mu |\mathbf{H}|^2) \mathbf{I} \right] , \quad (15)$$

where $\mathbf{I}$ is the identity tensor. The Maxwell-stress-tensor approach is broadly applicable and can be used to calculate optical forces on particles of arbitrary size and complex geometry. In practice, however, it generally requires numerical evaluation of the total field, including both the incident and scattered components, over the enclosing surface. Consequently, this formulation does not always provide direct physical insight into how the optical force is related to the intensity gradient, phase gradient, polarization structure, and orbital or spin momentum densities of the incident field⁴². For a small particle satisfying the Rayleigh condition $ka \ll 1$, the optical response can be described within the induced-dipole approximation⁴. The electric and magnetic

fields induce an electric dipole moment $\mathbf{d}_e$ and a magnetic dipole moment $\mathbf{d}_m$, respectively: $\mathbf{d}_e = \alpha_e \mathbf{E}$, $\mathbf{d}_m = \alpha_m \mu \mathbf{H}$, where α_e and α_m are the electric and magnetic polarizabilities of the particle. Within this approximation, the optical force can be expressed analytically, allowing a more transparent connection to be established between local field quantities and the resulting mechanical action on the particle. Conventional dipole-force theories were initially developed primarily for electric-dipole responses^{54,71–73}, but were subsequently extended to particles exhibiting both electric and magnetic dipole responses^{60,74–75}. The total force can then be decomposed as^{55–56,60}

$$\mathbf{F} = \mathbf{F}_e + \mathbf{F}_m + \mathbf{F}_{e-m} , \quad (16)$$

where $\mathbf{F}_e$ and $\mathbf{F}_m$ are the electric- and magnetic-dipole forces, respectively, and $\mathbf{F}_{e-m}$ is the additional force arising from interference between the radiation fields of the electric and magnetic dipoles. The electric-dipole force can first be written as

$$\mathbf{F}_e = \frac{1}{4} \text{Re} (\alpha_e) \nabla |\mathbf{E}|^2 + \frac{1}{2} \text{Im} (\alpha_e) \left\{ k \sqrt{\frac{\mu}{\varepsilon}} \text{Re} (\mathbf{E}^* \times \mathbf{H}) + \text{Im} [(\mathbf{E}^* \cdot \nabla) \mathbf{E}] \right\} . \quad (17)$$

The first term in Eq. (17) is associated with the electric-field energy-density gradient, whereas the remaining terms describe momentum transfer from the optical field to the particle and generally constitute a nonconservative radiation-pressure contribution. Using the momentum decomposition introduced above, Eq. (17) can be rearranged as

$$\mathbf{F}_e = \frac{2\pi n_m a^3}{c_0} \left(\frac{n_r^2 - 1}{n_r^2 + 2} \right) \nabla I + \frac{2\omega}{\varepsilon} \text{Im} (\alpha_e) \mathbf{P}_O^e , \quad (18)$$

where a is the particle radius, I is the optical intensity, c_0 is the speed of light in vacuum, and $n_r = n_p/n_m$ is the relative refractive index, with n_p and n_m denoting the refractive indices of the particle and surrounding medium, respectively. The first term in Eq. (18) is the well-known electric-field gradient force^{3,33,76}. For a dielectric particle with a refractive index higher than that of the surrounding medium, this force generally points toward regions of higher intensity and provides the restoring force required for stable optical trapping¹⁸. For a low-index particle, the force instead points toward intensity minima⁷⁷. Equation (18) therefore shows that the electric-dipole force contains two principal contributions: a conservative intensity-gradient force and a nonconservative momentum-transfer force associated with the electric part of the orbital momentum density.

Analogously, the magnetic-dipole force can be expressed as

$$\mathbf{F}_m = \frac{\mu^2}{4} \text{Re} (\alpha_m) \nabla |\mathbf{H}|^2 + 2\omega\mu \text{Im} (\alpha_m) \mathbf{P}_O^m . \quad (19)$$

Equation (19) shows that the magnetic-dipole force is governed primarily by the magnetic-field intensity gradient and the magnetic contribution to the orbital momentum

density. For most conventional dielectric particles, the magnetic response is weak at visible frequencies, and the magnetic-dipole contribution can usually be neglected. For high-refractive-index dielectric nanoparticles or magnetodielectric particles, however, the magnetic-dipole response and its coupling to the electric-dipole response may substantially modify the optical-force distribution.

When both electric and magnetic dipoles are simultaneously induced, interference between their radiation fields produces an additional electric-magnetic interaction force:

$$\begin{aligned}\mathbf{F}_{e-m} &= -\frac{k^4}{12\pi} \sqrt{\frac{\mu}{\varepsilon}} \operatorname{Re}(\mathbf{d}_e \times \mathbf{d}_m^*) \\ &= -\frac{\mu ck^4}{6\pi\varepsilon} \operatorname{Re}(\alpha_e \alpha_m^*) (\mathbf{P}_O + \mathbf{P}_S) + \frac{\mu ck^4}{6\pi\varepsilon} \operatorname{Im}(\alpha_e \alpha_m^*) \mathbf{P}_I.\end{aligned}\quad (20)$$

Equations (18)–(20) establish a direct connection between the local properties of the optical field and the time-averaged translational force acting on a dipolar particle. The isolated electric- and magnetic-dipole forces each contain an energy-density-gradient term and an orbital-momentum term, whereas the spin and imaginary momentum densities enter the translational force through electric-magnetic dipole interference. To make this field-to-force correspondence explicit, the total dipolar force in Eq. (16) can be reorganized according to the physical origins of its constituent terms as

$$\mathbf{F} = \mathbf{F}^{(G)} + \mathbf{F}^{(O)} + \mathbf{F}^{(S)} + \mathbf{F}^{(I)}. \quad (21)$$

The energy-density-gradient contribution is

$$\mathbf{F}^{(G)} = \frac{1}{4} \operatorname{Re}(\alpha_e) \nabla |\mathbf{E}|^2 + \frac{\mu^2}{4} \operatorname{Re}(\alpha_m) \nabla |\mathbf{H}|^2. \quad (22)$$

The orbital momentum density contribution is

$$\begin{aligned}\mathbf{F}^{(O)} &= \frac{2\omega}{\varepsilon} \operatorname{Im}(\alpha_e) \mathbf{P}_O^e + 2\omega\mu \operatorname{Im}(\alpha_m) \mathbf{P}_O^m \\ &\quad - \frac{\mu ck^4}{6\pi\varepsilon} \operatorname{Re}(\alpha_e \alpha_m^*) \mathbf{P}_O.\end{aligned}\quad (23)$$

The force induced by the spin momentum densities is

$$\mathbf{F}^{(S)} = -\frac{\mu ck^4}{6\pi\varepsilon} \operatorname{Re}(\alpha_e \alpha_m^*) \mathbf{P}_S, \quad (24)$$

and the force induced by the imaginary momentum

density is

$$\mathbf{F}^{(I)} = \frac{\mu ck^4}{6\pi\varepsilon} \operatorname{Im}(\alpha_e \alpha_m^*) \mathbf{P}_I. \quad (25)$$

Equations (22)–(25) provide a unified field-to-force correspondence within the electric-magnetic dipole approximation. The isolated electric- and magnetic-dipole forces each contain an energy-density-gradient term and an orbital-momentum term. The former is generally conservative and provides the restoring force in optical traps, whereas the latter describes nonconservative momentum transfer associated with absorption, scattering, local phase gradients, and radiation recoil.

By contrast, neither the spin momentum nor the imaginary Poynting momentum produces an independent time-averaged translational force through a single isotropic electric- or magnetic-dipole channel. At the electric-magnetic dipole level, $\mathbf{F}^{(S)}$ and $\mathbf{F}^{(I)}$ arise from interference between the induced electric and magnetic dipoles and are weighted by $\operatorname{Re}(\alpha_e \alpha_m^*)$ and $\operatorname{Im}(\alpha_e \alpha_m^*)$, respectively. The physical origins, particle-response requirements, and representative mechanical effects of the different force components are summarized in Table 1.

Beyond the dipole approximation, the forces associated with the spin momentum density $\mathbf{P}_S$ and imaginary Poynting momentum density $\mathbf{P}_I$ can acquire additional contributions through interference between higher-order multipoles^{55,56}. The Rayleigh scaling summarized in Table 1 is therefore restricted to the small-particle regime. For larger Mie particles, the total force can be evaluated using the Maxwell stress tensor, while multipolar decomposition helps identify the contributions associated with different optical-momentum components.

3 Generation and Control of Spatially Structured Light

3.1 Representative scalar and vector structured-light modes

In Cartesian coordinates, the Hermite-Gaussian (HG) modes form a set of orthogonal solutions to the paraxial wave equation. After omission of a common propagation

Table 1 | Physical origins, response requirements, and mechanical effects of different optical-force components.

Force component	Field quantity	Dipolar channel	Typical Rayleigh scaling	Mechanical effect
$\mathbf{F}^{(G)}$	$\nabla \mathbf{E} ^2, \nabla \mathbf{H} ^2$	$\mathbf{F}_e, \mathbf{F}_m$	$\propto a^3$	Conservative trapping
$\mathbf{F}^{(O)}$	$\mathbf{P}_O^e, \mathbf{P}_O^m, \mathbf{P}_O$	$\mathbf{F}_e, \mathbf{F}_m, \mathbf{F}_{e-m}$	Absorptive: $\propto a^3$; nonabsorptive scattering: $\propto a^6$	Radiation pressure and directed transport
$\mathbf{F}^{(S)}$	$\mathbf{P}_S$	$\mathbf{F}_{e-m}$, weighted by $\operatorname{Re}(\alpha_e \alpha_m^*)$	Electric-magnetic dipoles: $\propto a^6$; Isolated electric or magnetic dipole: 0	Spin-momentum-dependent lateral force
$\mathbf{F}^{(I)}$	$\mathbf{P}_I$	$\mathbf{F}_{e-m}$, weighted by $\operatorname{Im}(\alpha_e \alpha_m^*)$	Response-dependent; Isolated electric or magnetic dipole: 0	Imaginary-momentum-mediated recoil force

factor, their complex amplitude can be written as^{78–79}

$$U_{l,m}^{HG}(x, y, z) = A_{l,m} \frac{\omega_0}{\omega(z)} H_l \left(\frac{\sqrt{2}x}{\omega(z)} \right) H_m \left(\frac{\sqrt{2}y}{\omega(z)} \right) \times \exp \left[-\frac{x^2 + y^2}{\omega^2(z)} \right] \exp \left[-ik \frac{x^2 + y^2}{2R(z)} - i(l + m + 1) \tan^{-1} \left(\frac{z}{z_0} \right) \right], \quad (26)$$

where $A_{l,m}$ is a normalization constant, $\omega(z) = \omega_0 \sqrt{1 + (z/z_0)^2}$ is the beam radius, ω_0 is the waist radius, $z_0 = \pi \omega_0^2 / \lambda$ is the Rayleigh range, and $R(z) = z[1 + (z_0/z)^2]$ is the radius of curvature of the wavefront. H_l and H_m are Hermite polynomials of orders l and m , respectively.

The left-hand side of Fig. 3 shows nine HG modes of different orders. As the transverse indices l and m increase, the intensity profile divides into multiple lobes along the x and yy directions, respectively. Adjacent lobes generally differ in phase by π . Both the intensity pattern and phase discontinuities of an HG mode can therefore be controlled directly through the two mode indices. For optical manipulation, such Cartesian multilobe structures provide multiple spatially separated intensity gradients and anisotropic trapping landscapes, making HG modes suitable for multipoint confinement and directional organization of particles⁴⁰.

In cylindrical coordinates, the corresponding orthogonal solutions are the Laguerre–Gaussian (LG) modes^{79–80}:

$$U_{m,p}^{LG}(\rho, \varphi, z) = A_{m,p} \frac{\omega_0}{\omega(z)} \left(\frac{\sqrt{2}\rho}{\omega(z)} \right)^{|m|} \exp \left(-\frac{\rho^2}{\omega^2(z)} \right) L_p^{|m|} \left(\frac{2\rho^2}{\omega^2(z)} \right) \times \exp \left[-ik \frac{\rho^2}{2R(z)} + im\varphi - i(2p + |m| + 1) \tan^{-1} \left(\frac{z}{z_0} \right) \right], \quad (27)$$

where $A_{m,p}$ is an amplitude coefficient, $L_p^{|m|}$ is an associated Laguerre polynomial, and m and p are the azimuthal and radial indices, respectively.

As shown on the right-hand side of Fig. 3, increasing the radial index p increases the number of radial intensity rings. When the azimuthal index satisfies $m \neq 0$, the beam carries a vortex phase $\exp(im\varphi)$, whose phase accumulates by $2\pi m$ over one complete revolution around the beam axis. Under the paraxial approximation, each photon carries an orbital angular momentum of $m\hbar$. Because a phase singularity lies on the axis, an LG mode with $m \neq 0$ generally possesses a central intensity null. These features have direct mechanical consequences in optical manipulation: the annular intensity distribution provides radial confinement, whereas the helical phase generates azimuthal momentum flow and enables orbital-angular-momentum transfer, making LG vortex beams particularly suitable for ring-shaped trapping, orbital

transport, and particle rotation^{43–45,73,81–82}.

Another class of rotationally symmetric solutions in cylindrical coordinates is the Bessel beam⁴⁷:

$$U_m^B(\rho, \varphi, z) = A_m J_m(k_\rho \rho) \exp(im\varphi) \exp(ik_z z), \quad (28)$$

where A_m is an amplitude coefficient, J_m is the Bessel function of order m , and k_ρ and k_z are the transverse and longitudinal wavevector components, satisfying $k^2 = k_\rho^2 + k_z^2$. Figure 4 shows the complex-amplitude distributions of Bessel modes with $m = 0$, $m = 4$, and $m = 8$. As shown in Fig. 4(a), the zeroth-order Bessel beam has a bright central lobe. At higher orders, a helical phase and a central intensity null appear, and the phase varies more rapidly along the azimuthal direction, as illustrated in Fig. 4(b, c). Both the extended depth of field and self-healing behavior provide robust optical confinement over an extended propagation range, while the axially elongated multi-ring structure enables simultaneous trapping at multiple axial positions. These characteristics make Bessel beams particularly suitable for multiplane trapping and long-range axial transport^{46–47,83}.

In addition to the representative scalar modes described above, the paraxial wave equation has other analytical solutions under different coordinate systems or boundary conditions, including Bessel–Gaussian modes^{84–85}, Ince–Gaussian modes⁸⁶, and Mathieu beams^{87–88}. Their distinct symmetries, propagation stability, and sidelobe structures provide a broad range of optical-potential landscapes and transport trajectories for optical micromanipulation⁴⁰.

Spatially varying polarization introduces an additional degree of freedom. In conventional linearly, circularly, or elliptically polarized beams, the polarization state is approximately uniform across the beam cross section. When the polarization direction or ellipticity varies with transverse position, the field becomes vectorial and is generally referred to as a vector beam. Radially and azimuthally polarized beams are two representative cylindrical vector beams, with local polarization directions oriented along the radial and azimuthal directions, respectively.

Both can be constructed by superposing orthogonal HG modes carrying orthogonal polarizations. A radially polarized field $\mathbf{E}_r$ and an azimuthally polarized field $\mathbf{E}_\varphi$ may be written as^{51,89}

$$\mathbf{E}_r = \text{HG}_{1,0} \hat{\mathbf{e}}_x + \text{HG}_{0,1} \hat{\mathbf{e}}_y, \quad (29)$$

$$\mathbf{E}_\varphi = -\text{HG}_{0,1} \hat{\mathbf{e}}_x + \text{HG}_{1,0} \hat{\mathbf{e}}_y, \quad (30)$$

where $\hat{\mathbf{e}}_x$ and $\hat{\mathbf{e}}_y$ are mutually orthogonal linear-polarization basis vectors.

As shown in Fig. 5, the electric-field vectors of radially polarized light are oriented along the radial direction, whereas those of azimuthally polarized light are oriented tangentially. Owing to their distinctive spatially inhomogeneous polarization distributions, these vector beams have found numerous applications in high-resolution imaging,

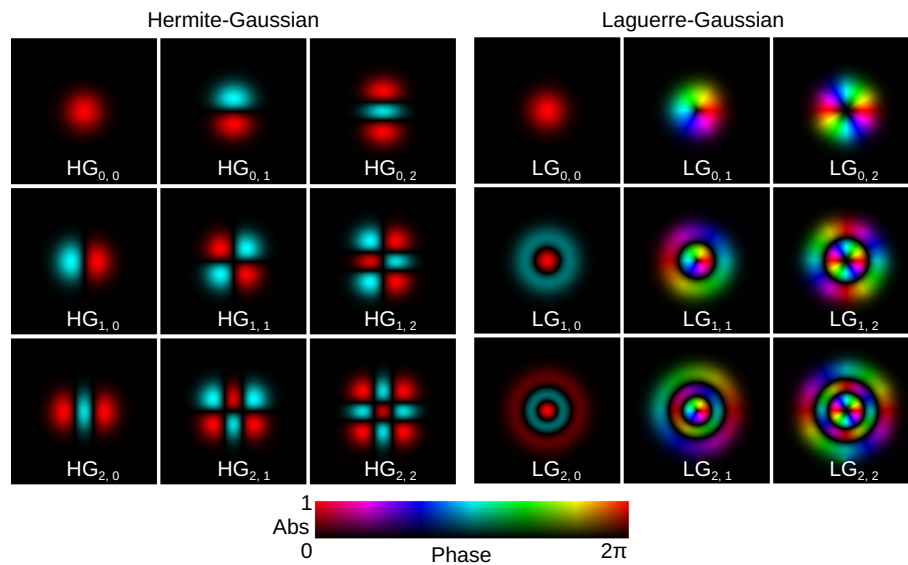


Fig. 3 | Intensity and phase distributions of Hermite–Gaussian and Laguerre–Gaussian modes.

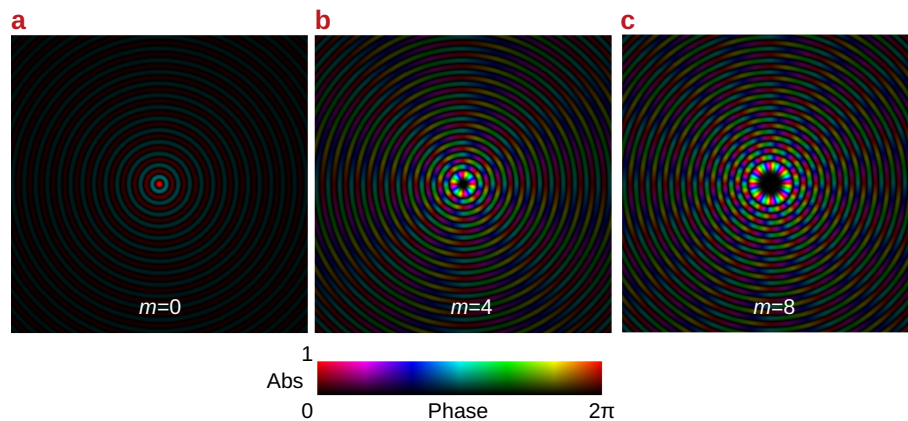


Fig. 4 | Complex-amplitude distributions of Bessel beams of different orders. (a) $m = 0$; (b) $m = 4$; and (c) $m = 8$. The color represents phase. The zeroth-order mode has an on-axis bright spot, whereas higher-order modes exhibit a central dark core and an azimuthally varying helical phase.

plasmonic focusing, and particle manipulation⁵¹. For optical manipulation in particular, spatially varying polarization provides an additional degree of freedom for tailoring local spin density, polarization gradients, and polarization-dependent optical forces and torques, complementing the intensity- and phase-based mechanisms of scalar structured fields.

3.2 Two-dimensional structured-field control

The generation of an arbitrary complex-amplitude field requires the target amplitude and phase constraints to be satisfied simultaneously. When both quantities must be independently prescribed in the output plane, a simple analytical solution for a phase-only hologram is generally unavailable. The conventional Gerchberg–Saxton (GS) algorithm iteratively propagates between the input and output planes and alternately imposes the relevant constraints⁹⁰. It

is effective for reconstructing a target intensity distribution, but for an arbitrary complex field, the amplitude and phase constraints usually cannot both be satisfied exactly⁹¹.

To address this problem, researchers introduced two completely overlapping virtual planes at the output, one imposing amplitude constraints and the other imposing phase constraints, with the constrained regions chosen to be complementary. This arrangement enables cooperative optimization of the complex-amplitude distribution over the entire output plane^{92–93}. As shown in Fig. 6(a), the method can simultaneously reconstruct a prescribed intensity and phase distribution along the same curve, allowing a continuous bright trajectory to carry a designed local phase gradient.

In order to improve modulation efficiency while maintaining trajectory following and simultaneously controlling amplitude and phase, phase-only shaping methods such as

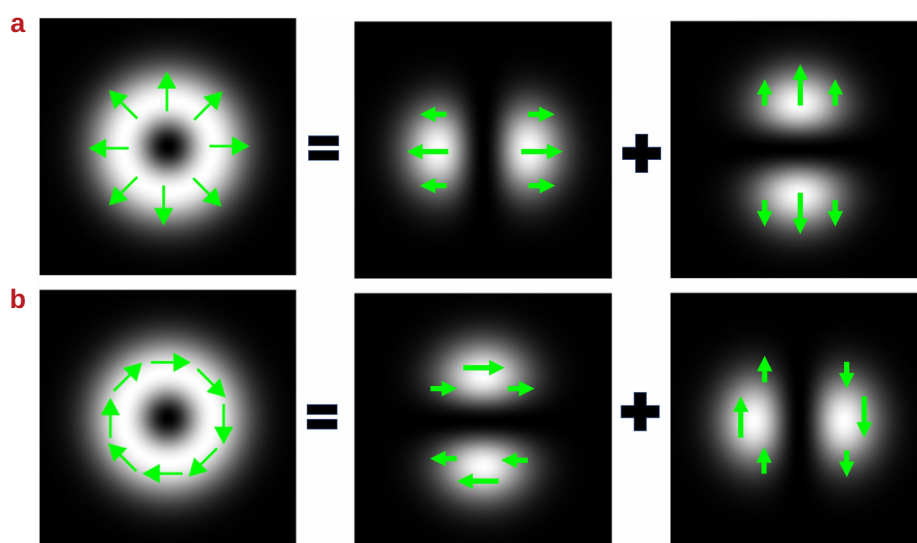


Fig. 5 | Polarization distributions of two representative cylindrical vector beams. (a) Radially polarized beam and (b) azimuthally polarized beam. Both fields can be formed by superposing orthogonally polarized HG modes. The arrows indicate the local linear-polarization direction.

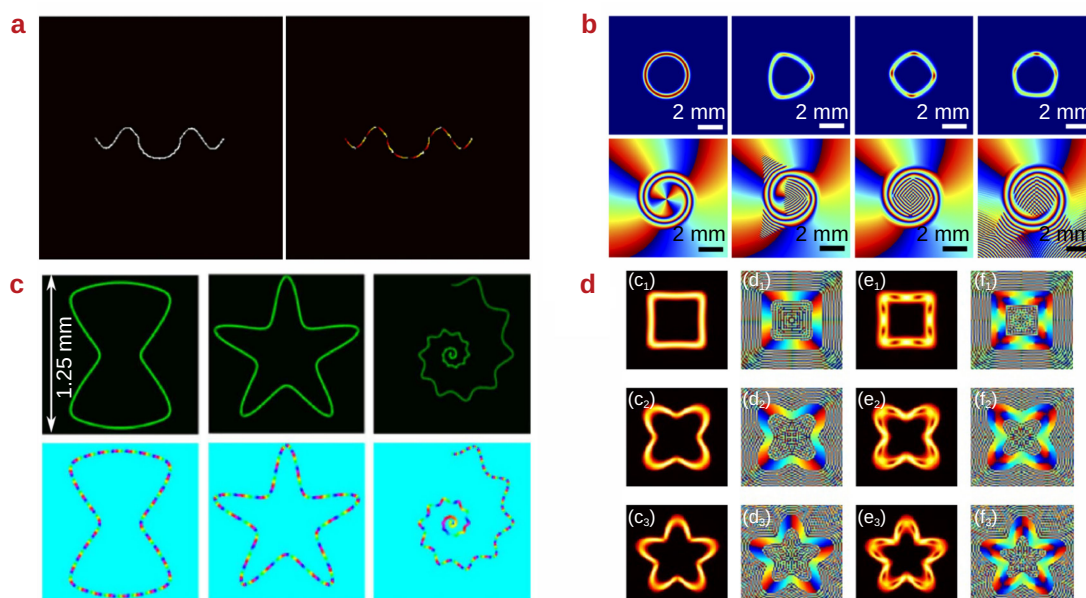


Fig. 6 | Two-dimensional joint control of arbitrary intensity and phase distributions. (a) Iterative algorithm for simultaneous amplitude and phase constraints⁹³; (b) phase-only wavefront shaping⁹⁸; (c) curve-structured method for generating continuous intensity trajectories and phase distributions¹⁰⁰; (d) curve multiplication and vortex phase array design¹⁰⁵. Figure reproduced from: (a) ref.⁹³, Springer Nature; (b) ref.⁹⁸, Optica Publishing Group; (c) refs.¹⁰⁰, Springer Nature; (d) ref.¹⁰⁵, Optica Publishing Group.

the freeform-lens method^{94–97} and cross-phase modulation⁹⁸ have also been developed. As shown in Fig. 6(b), these methods can generate optical fields with predetermined geometrical shapes and phase structures. In experimental systems employing phase-only devices, these phase-only approaches achieve high efficiency, but in principle, the quality and precision of the generated fields are still inferior to those produced by complex-amplitude modulation.

On another technical track, significant progress has also been made in the study of curve-structured optical beams. In

2004, Abramochkin et al. theoretically proposed a class of Gaussian-distributed structured beams whose intensity structure remains invariant during focusing⁹⁹. Subsequently, Rodrigo et al. further developed this approach^{100–103}, obtaining curve-structured optical fields with steep intensity gradients by removing the Gaussian term, and found that the incident fields of such beams can be directly obtained analytically via parametric equations. As shown in Fig. 6(c), by incorporating the corresponding parametric equations, one can generate hourglass-shaped, star-shaped, spiral, and

other continuous intensity trajectories, while independently prescribing the phase distribution along each curve. This method is therefore particularly suitable for particle transport along curved paths. In recent years, this technique has been extended to corresponding flat-top beams¹⁰⁴, and by partially superimposing two flat-top beams with different radii, it has enabled the manipulation of phase vortex arrays distributed along various intensity trajectories¹⁰⁵, as shown in Fig. 6(d).

3.3 Three-dimensional structured-field control

Three-dimensional optical field manipulation requires not only reconstructing the target distribution in a transverse plane, but also controlling the evolution along the propagation direction. Early methods mainly relied on multi-plane holographic superposition and axial multifocal scanning^{106–109}, which are suitable for constructing discrete lattices but are difficult to generate continuous, smooth three-dimensional trajectories with independent phase

control. The 3D GS algorithm can iteratively optimize across multiple target planes to reconstruct complex structures such as helices, crossed lines, and volume lattices¹¹⁰, as shown in Fig. 7(a); while inverse design based on the wave equation directly solves for the incident wavefront from the target three-dimensional field¹¹¹, as shown in Fig. 7(b). These methods greatly expand the utilization of spatial degrees of freedom, enabling the fabrication of controllable chiral microstructures inside isotropic polymers using three-dimensional structured light fields¹¹², and extending the dimension of particle manipulation from two-dimensional planes to three-dimensional space³⁴.

True three-dimensional optical trapping also requires sufficiently strong axial confinement to overcome the scattering force along the propagation direction. As shown in Fig. 7(c), the Grier group first proposed curve-structured optical fields with extremely high axial intensity gradients by solving the inverse diffraction problem^{113,114}, and experimentally verified their stable three-dimensional trapping capability. Subsequently, the Rodrigo group developed an analyti-

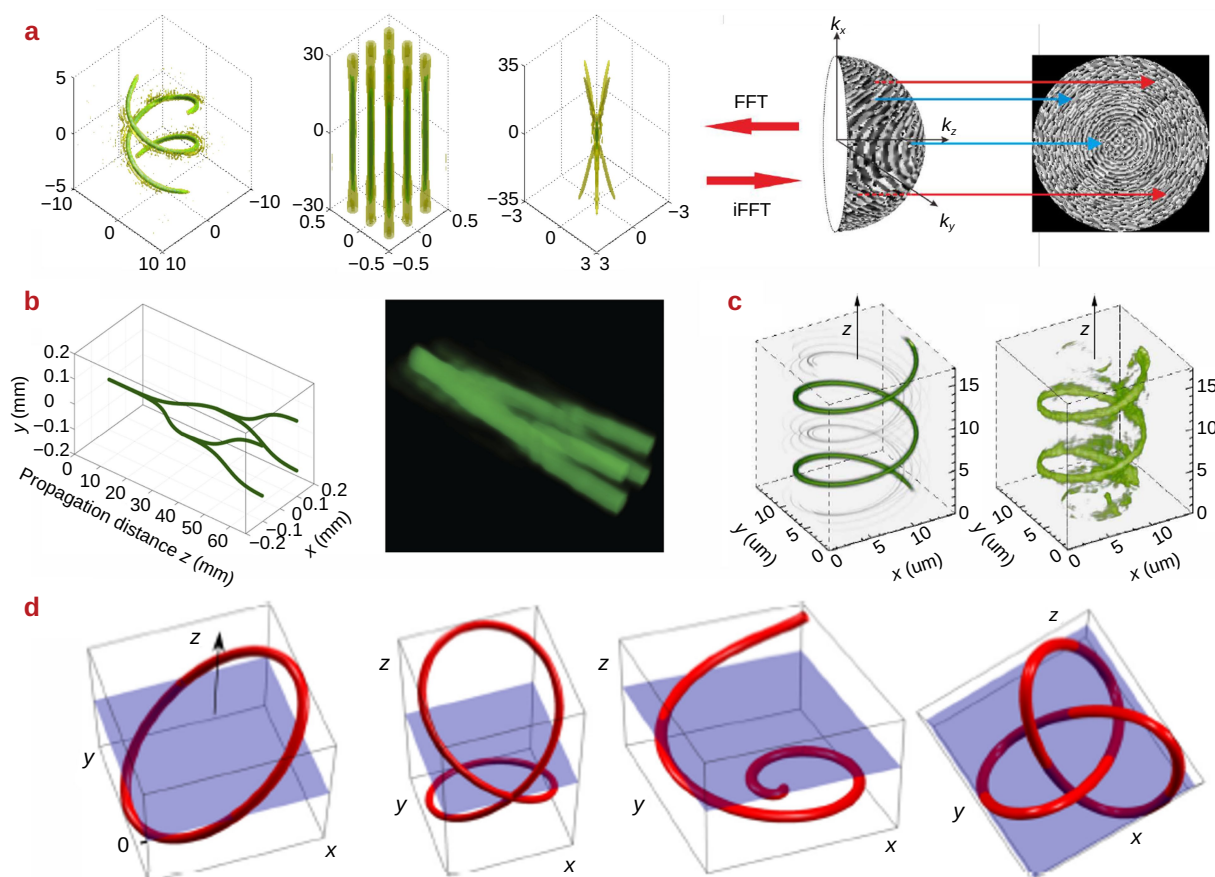


Fig. 7 | Representative methods for constructing three-dimensional scalar optical fields. (a) Three-dimensional GS iteration for generating double helices, line arrays, and crossed structures¹¹⁰; (b) inverse design of a three-dimensional field based on the wave equation¹¹¹; (c) solenoid beams¹¹³; and (d) three-dimensional curved focal fields of different geometries constructed using an axial-displacement phase³⁴. Figure reproduced from: (a) ref.¹¹⁰, IOP Publishing; (b) ref.¹¹¹, Optica Publishing Group; (c) ref.¹¹³, Optica Publishing Group; (d) ref.³⁴, Optica Publishing Group.

cal method by introducing an axial shift factor, which allows direct calculation of the corresponding incident field once the parametric equations of the desired three-dimensional curve are specified. As illustrated in Fig. 7(d), this method can generate tilted rings, spatial closed curves, Archimedean spirals, and clover-shaped focal lines. For a three-dimensional focal curve described by $c(t) = (x_0(t), y_0(t), z_0(t))$, the input field can be written as^{34,65}

$$H(x, y) = \int_0^T g(t) \gamma(x, y, t) \xi(x, y, t) dt, \quad (31)$$

where

$$\begin{cases} g(t) = \eta(t) \exp(i\psi(t)) \sqrt{x_0'(t)^2 + y_0'(t)^2 + z_0'(t)^2} \\ \gamma(x, y, t) = \exp\left(-\frac{i2\pi}{\lambda f} (xy_0(t) + yx_0(t))\right) \\ \xi(x, y, t) = \exp\left(i\pi \frac{x^2 + y^2}{\lambda f^2} z_0(t)\right). \end{cases} \quad (32)$$

Here, $\eta(t)$ controls the amplitude variation along the

curve, whereas $\psi(t)$ determines the local phase gradient along the curve. The factor $\gamma(x, y, t)$ specifies the transverse position of each focal point, enabling consecutive foci to trace out a curve like a brush painting dots into lines. The quadratic phase factor $\xi(x, y, t)$ controls the defocus distance of each point relative to the reference focal plane. The parameters f and λ denote the focal length of the focusing lens and the wavelength of the incident light, respectively. This method can generate helices extending either upward or downward¹¹⁵ and can be combined with various natural-curve equations^{116–117} to construct more complex three-dimensional structures¹¹⁸.

In addition to directly designing three-dimensional focal lines, the evolution of fundamental modes during propagation can also be manipulated. Starting from the two-dimensional Fourier transform in cylindrical coordinates, Yan et al. proposed a scheme for self-scaling and self-accelerating manipulation of the main lobe of Bessel beams in three-dimensional space¹¹⁹. As shown in Fig. 8, the left column presents the corresponding energy flow lines, while the right

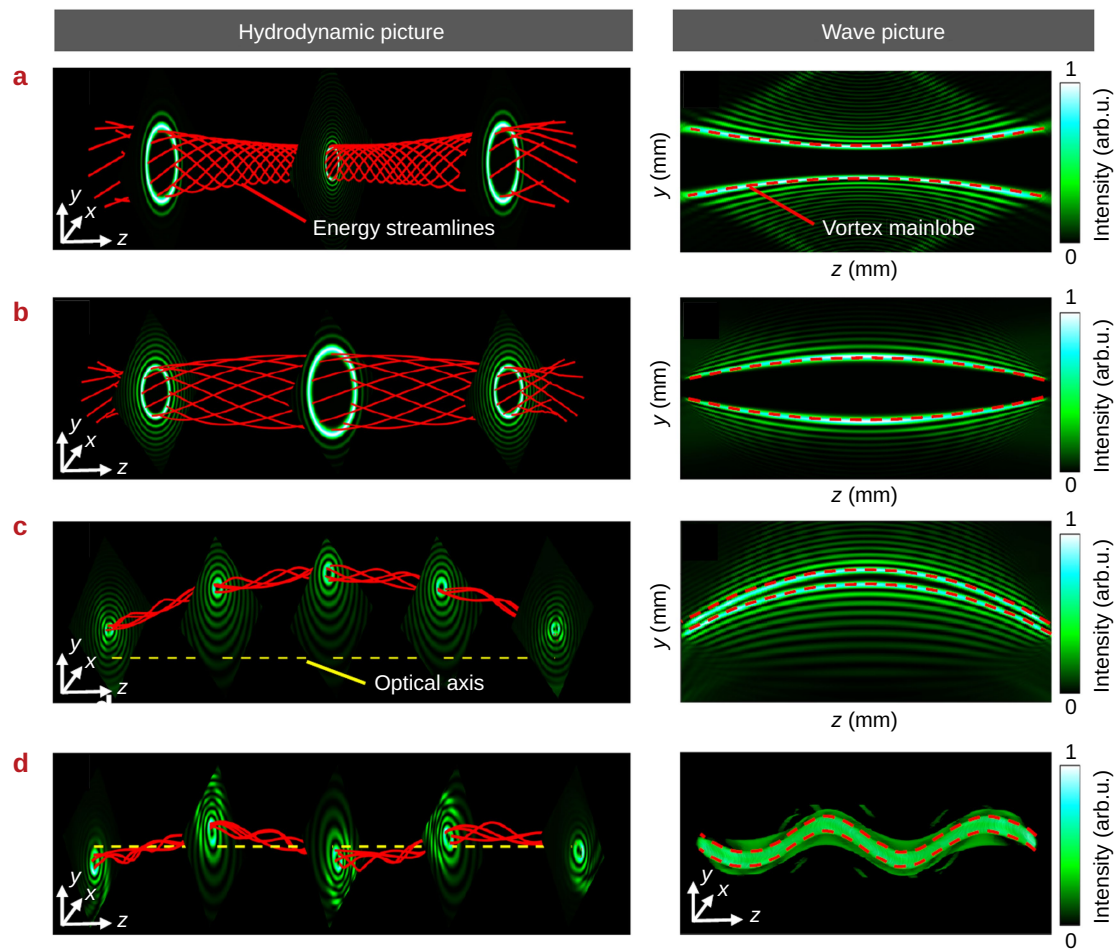


Fig. 8 | Control of the three-dimensional evolution of Bessel beams. The left-hand column shows the energy-flow-line distributions, and the right-hand column shows the corresponding intensity evolution, including axial propagation, self-scaling, curved self-acceleration, and periodic self-acceleration¹¹⁹. Figure reproduced from: ref. ¹¹⁹, Springer Nature.

column shows the evolution of the intensity main lobe with propagation distance. In Fig. 8(a, b), the energy flow lines contract or expand, enabling the Bessel main lobe to adaptively change its transverse size along the propagation direction; in Fig. 8(c, d), the main lobe propagates along curved or periodically oscillating trajectories, respectively, exhibiting controllable self-accelerating behavior. The correspondence between the energy flow lines and the intensity trajectories intuitively reveals the dynamical process of the three-dimensional beam evolution.

The corresponding angular spectrum can be expressed as¹¹⁹

$$\begin{cases} E_i(\sqrt{k^2 - k_z^2}, \varphi) = \frac{1}{\text{rect}(k_z/2k)} e^{im\varphi} \int_{-\infty}^{\infty} e^{i\varphi(z)} e^{ik_z z} e^{-ik_z z} dz \\ \varphi(z) = k_x g(z) + k_y h(z) \end{cases} \quad (33)$$

A longitudinal spatial-frequency-shift phase varying with propagation distance z is introduced. The functions $g(z)$ and $h(z)$ control the horizontal and vertical displacements of the Bessel-beam main lobe from the optical axis at the propagation distance z . In combination with Eq. (5), the spatial evolution of the optical energy-flow lines can also be calculated. Notably, the scheme can preserve spatial characteristics such as the radius of the main vortex ring while controlling the three-dimensional trajectory, producing a so-called “three-dimensional perfect vortex”¹²⁰. This concept extends the conventional two-dimensional perfect vortex into the propagation dimension^{121–122}.

A three-dimensional vector field additionally requires the polarization state to be prescribed at every point in space. A commonly used strategy decomposes the target vector field into two orthogonally polarized components, calculates the corresponding three-dimensional scalar holograms independently, and then reconstructs the complete vector field by polarization multiplexing and coaxial superposition^{123–124}. As shown in Fig. 9(a), separate control of the two polarization components enables the generation of focal fields with prescribed local linear-polarization directions at different axial positions.

By selecting different polarization bases and controlling the relative phase between the two components, the method can be extended to hybrid-order vector structured fields¹²⁵. Figure 9(b) represents different polarization states on the Poincaré sphere and illustrates the continuous evolution of the local polarization during propagation. Nevertheless, these approaches rely mainly on phase and polarization multiplexing and make limited independent use of the amplitude degree of freedom.

With amplitude control included, locally linearly polarized fields can be generalized to fully vectorial three-dimensional focal fields. As illustrated in Fig. 9(c), simultaneous control of amplitude, phase, and polarization makes it possible to establish, both theoretically and experimentally, a one-to-one correspondence between the spatial intensity distri-

bution and the distributions of the three Stokes parameters¹²⁶. The optical intensity, polarization ellipticity, and polarization orientation can therefore be independently prescribed at each point of the target three-dimensional field.

Beyond conventional spatial-light-modulation approaches, structural functional materials, particularly metasurfaces, provide a compact platform for vector-field engineering by locally tailoring amplitude, phase, and polarization responses at the subwavelength scale¹²⁷. By combining geometric and dynamic phases, spin-decoupled metasurfaces can further realize independent transformations of different polarization and orbital-angular-momentum channels, illustrating the potential of meta-optics for compact generation, transformation, and characterization of high-dimensional structured fields¹²⁸.

For classical scalar modes such as Bessel or Gaussian-type beams, the continuous evolution of the polarization state with propagation distance can also be achieved by adjusting the amplitude ratio and the longitudinal propagation phase difference between the two orthogonal polarization bases^{129–130}. As shown in Fig. 10, the two orthogonal polarization components accumulate different phases during propagation, causing the local polarization state to vary periodically along the optical axis and exhibiting non-separable longitudinal oscillatory variations¹³⁰. This provides an important foundation for constructing propagation-stable high-dimensional vector optical fields and for enabling multi-degree-of-freedom optical manipulation.

4 Applications of Structured Light in Optical Tweezers

The applications of structured light are organized below according to the optical degrees of freedom involved and the dimensionality of particle manipulation. We first consider scalar structured fields, progressing from spatially tailored trapping landscapes and planar intensity–phase control to three-dimensional and axial manipulation. We then introduce polarization as an additional degree of freedom, ranging from spin-mediated rotation and translation to mechanical effects arising from spatially varying vector fields. Finally, light-driven micro- and nanorobots are discussed as an emerging direction in which the optical response of the manipulated object itself is engineered to achieve increasingly complex mechanical functions.

4.1 Scalar Structured-Light Manipulation

Scalar structured fields primarily utilize spatially tailored intensity and phase distributions with uniform polarization to reshape trapping potentials and local momentum flow, with applications spanning multipoint confinement, planar transport, dark trapping, and three-dimensional manipula-

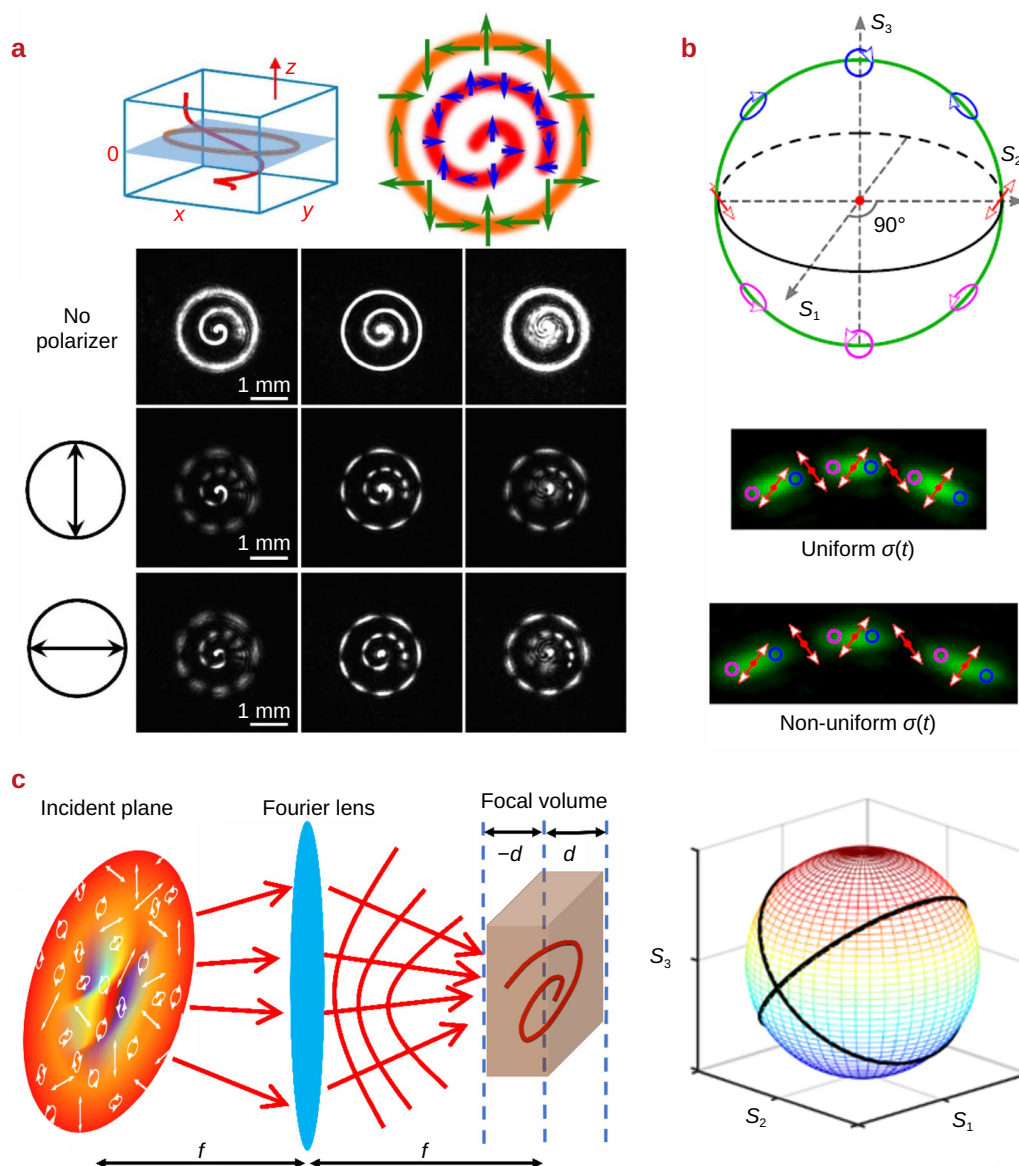


Fig. 9 | Generation of three-dimensional vector structured fields. (a) Decomposition of the target field into orthogonally polarized components for reconstructing a three-dimensional locally linearly polarized focal field¹²³; (b) generation of hybrid vector structured fields through relative-phase control and polarization-basis transformation¹²⁵; and (c) construction of general three-dimensional fully vectorial focal fields through joint control of amplitude, phase, and polarization¹²⁶. Figure reproduced from: (a) ref.¹²³, Optica Publishing Group; (b) ref.¹²⁵, Springer Nature; (c) ref.¹²⁶, Chinese Laser Press.

tion.

4.1.1 Multipoint Trapping and Array Manipulation

A conventional single-focus Gaussian optical trap, however, provides only one localized trapping site and is therefore poorly suited to the parallel manipulation, complex spatial organization, and dynamic assembly of multiple particles. Holographic optical tweezers overcome this limitation by transforming a single incident beam into multiple independently addressable foci, enabling flexible control over the number, position, and three-dimensional arrangement of

optical traps.

Holographic optical tweezers can arrange multiple particles along the same projected line while positioning them at different axial planes spanning from $-5\ \mu\text{m}$ to $5\ \mu\text{m}$, thereby enabling genuinely three-dimensional particle organization³¹. Beyond regular arrays, dynamically controlled traps can translate particles sequentially along prescribed paths, allowing particle patterns to be directly “written” and providing a flexible approach to particle assembly and pattern-based microfabrication¹³¹.

Planar optical devices provide a complementary route

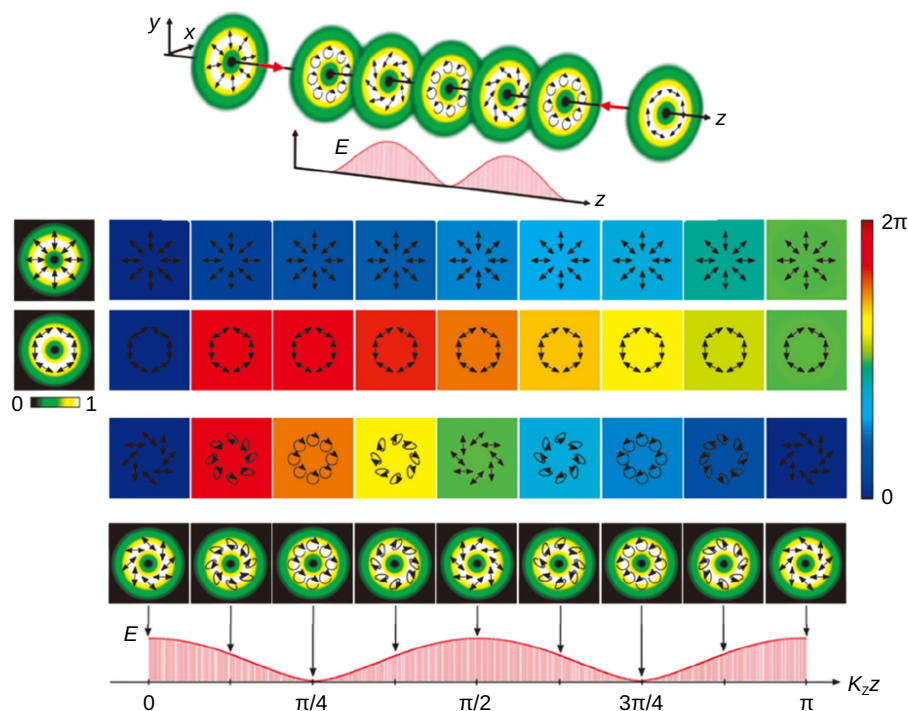


Fig. 10 | Propagation evolution of a three-dimensional vector field and characterization of spatial-polarization nonseparability¹³⁰. Figure reproduced from: ref.¹³⁰, Springer Nature.

toward larger and more compact optical-trap arrays. By engineering subwavelength meta-atoms, holographic meta-surfaces can generate diverse focal-array geometries with high optical efficiency and support the simultaneous trapping of large numbers of atoms or particles¹³². Recent meta-surface platforms have experimentally demonstrated more than 250 single-atom optical tweezers and, owing to their submicrometer pixel sizes and high pixel densities, provide a route toward arrays containing more than 10^5 trapping sites¹³², offering considerable potential for neutral-atom quantum simulation and quantum computing¹³³. Wavelength and polarization multiplexing can further be incorporated to switch focal positions between different orthogonal planes, enabling three-dimensional particle transport¹³⁴. These monolithic devices reduce the complexity of free-space optical arrangements and provide a promising platform for compact, efficient, and highly integrated optical manipulation.

In addition to externally prescribed trap arrays, multiple particles can self-organize through optical binding, in which light scattered by one particle modifies the optical forces acting on neighboring particles, giving rise to light-mediated interparticle interactions and stable collective configurations^{17,19}. Structured light provides additional degrees of freedom for tailoring these interactions. For example, counterpropagating Bessel beams can support long-range longitudinal optical binding and particle-chain formation¹³⁵, whereas the spatial intensity and polarization structures of

Laguerre–Gaussian beams can modify the optical-binding potential and equilibrium configurations¹³⁶. Transverse phase gradients can further regulate the interaction landscape and dynamics of optically bound nanoparticles¹³⁷. Optical binding provides a complementary route to multi-particle manipulation, extending structured-light control from prescribed trapping arrays to collective self-organization.

4.1.2 Joint Intensity and Phase Control for Planar Transport

For planar transport with scalar structured fields, intensity and phase generally perform complementary mechanical functions. The intensity distribution forms the transverse trapping potential, whereas the phase distribution reshapes the local phase gradient and momentum flow to propel particles along a prescribed trajectory. Joint control of these two degrees of freedom greatly expands the range of accessible particle motions.

An optical vortex provides a representative example of such intensity–phase cooperation. As shown in Fig. 11(a), the annular intensity distribution confines particles whose refractive index is higher than that of the surrounding medium near the bright ring, whereas the helical phase generates azimuthal momentum flow along the ring. Through the transfer of orbital angular momentum, trapped particles can continuously orbit the optical axis. Their orbital direction and speed can be controlled through the

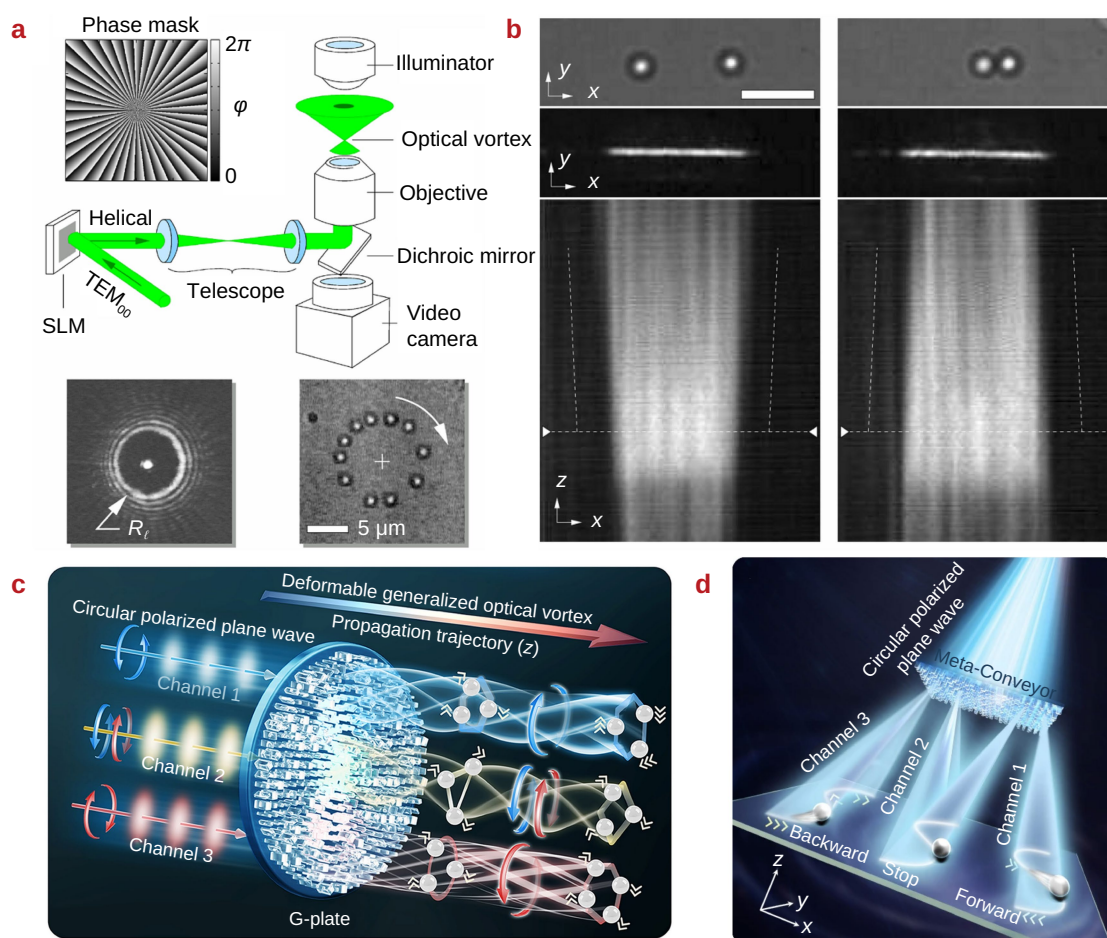


Fig. 11 | Joint intensity and phase control of particle motion. (a) Orbital motion produced by annular intensity confinement and helical-phase driving⁴³; (b) convergence and separation of particles along a linear optical track³³; and (c, d) polarization-multiplexed metasurfaces generating closed and open trajectories for multifunctional particle transport^{138–139}. Figure reproduced from: (a) ref.⁴³, American Physical Society; (b) ref.³³, American Physical Society; (c) ref.¹³⁸, American Association for the Advancement of Science; (d) ref.¹³⁹, Springer Nature.

topological charge, optical power, and particle properties⁴³.

For open trajectories, the phase gradient can be designed independently along a bright intensity line. As illustrated in Fig. 11(b), when the phase gradients at both ends of the line point toward its center, particles initially located near the two ends move toward one another and converge. When the gradients point away from the center, the particles move in opposite directions and separate³³. This result demonstrates that the particle-transport state can be switched solely by modifying the phase distribution while leaving the intensity trajectory essentially unchanged.

More recently, joint intensity and phase control has been extended to metasurface platforms. By engineering the geometry and in-plane orientation of anisotropic meta-atoms, the phase retardation and polarization-conversion terms of the local Jones matrix can be controlled point by point. The two orthogonal polarization components of the incident field can consequently be manipulated independently. Figure 11(c, d) show closed and open trajectories

generated by polarization-multiplexed metasurfaces. Different incident polarization states are encoded into independent optical channels, allowing a single device to switch between multiple transport modes and thereby enabling efficient, multifunctional particle conveyance^{138,139}.

4.1.3 Dark-Trap Manipulation of Low-Index and Absorbing Particles

Conventional optical tweezers use a tightly focused field to produce a strong intensity gradient that attracts particles with a refractive index higher than that of the surrounding medium toward high-intensity regions, typically near the beam focus¹⁸. Such particles therefore exhibit a pronounced “light-seeking” response. By contrast, particles with a refractive index lower than that of the surrounding medium are repelled from high-intensity regions. Absorbing particles may additionally experience strong radiation pressure and photothermal effects, making stable confinement in a conventional bright focus difficult. To manipulate these

“dark-seeking” particles, a variety of structured-field and computer-generated-holography-based dark traps have been developed^{77,140–144}.

Early studies used the central dark core of a Gaussian vortex beam to trap low-index particles^{77,140–141}. Higher-order Bessel beams were subsequently employed to construct dark traps through their multiring intensity distributions. As shown in Fig. 12(a), a low-index particle can be confined in the dark region between two neighboring bright rings and guided along that region¹⁴². Another strategy is to interfere two optical beams to produce periodic bright and dark fringes, within which particles can be trapped and translated by shifting the fringe pattern¹⁴³.

The dimensions of these dark regions are generally determined by the beam parameters or interference period and therefore offer only limited tunability. When the particle size is poorly matched to the dark-region width, trapping stability is substantially reduced. To alleviate this problem, quasi-perfect vortex beams have been used to form annular dark cores with adjustable dimensions, enabling low-refractive-index particles to be trapped and rotated¹⁴⁴. Nevertheless, their rotationally symmetric geometry usually confines particle transport to circular trajectories and does not readily support arbitrary path control¹⁴⁵.

Figure 12(b) shows a difference-of-Gaussians dark trap. In this method, two Gaussian fields are subtracted to create a localized dark core, which can trap an individual low-index particle and transport it by moving the dark region¹⁴⁶. However, each change in particle position or motion state generally requires a newly calculated hologram. Continuous transport requires repeated hologram updates and is therefore limited by hologram-generation and spatial-light-modulator refresh rates. Speckle fields can also trap low-index particles in their numerous randomly distributed dark regions¹⁴⁷, but the disordered nature of the field makes it difficult to define deterministic and continuous transport routes.

To achieve controlled transport along prescribed dark trajectories, Yuan et al. proposed superposing two arbitrary curved beams carrying independently designed phase distributions¹⁴⁸. As illustrated in Fig. 12(c), the two bright curves form bright trapping channels for high-index particles, while the region between them creates a continuous dark channel for low-index particles. By independently designing the phase gradient and orbital momentum flow along the curved paths, high- and low-index particles can be driven along different trajectories. Experiments demonstrated their simultaneous trapping, transport, and routing. A similar

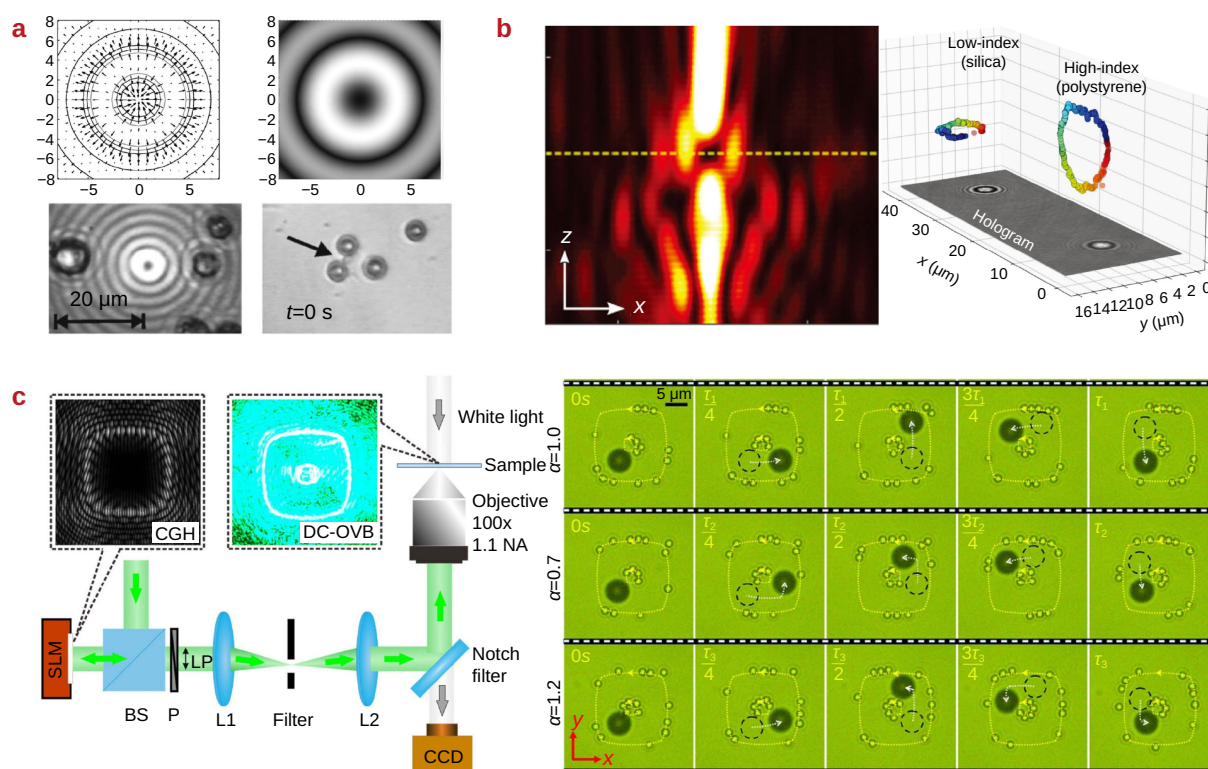


Fig. 12 | Dark-trap manipulation of low-refractive-index particles. (a) Trapping in the dark region between adjacent bright rings of a higher-order Bessel beam¹⁴²; (b) a difference-of-Gaussians dark trap¹⁴⁶; and (c) a controllable dark trajectory formed between two curved bright tracks for multipath manipulation of high- and low-index particles¹⁴⁸. Figure reproduced from: (a) ref.¹⁴², American Physical Society; (b) ref.¹⁴⁶, Optica Publishing Group; (c) ref.¹⁴⁸, Chinese Laser Press and Optica Publishing Group.

concept can also be implemented using a free-lens approach, providing a flexible platform for continuous transport of low-index particles⁹⁶.

Overall, conventional vortex, Bessel, and interference-based dark traps provide relatively simple and stable confinement but offer limited geometric flexibility, whereas holographically or computationally designed dark fields enable more tunable and complex trajectories at the cost of increased modulation complexity. The appropriate strategy therefore depends on the required balance between trapping stability and dynamic controllability.

For absorbing particles, structured light can also produce heat-mediated mechanical effects that differ fundamentally from direct electromagnetic gradient and radiation-pressure forces. In gaseous environments, optical absorption creates asymmetric particle heating and momentum exchange with surrounding gas molecules, giving rise to photophoretic forces. Optical vortex beams have been used to confine absorbing particles in their low-intensity cores and enable long-range transport over macroscopic distances¹⁴⁹. More remarkably, switching between azimuthally and radially polarized vector beams can reverse the photophoretic force on partially absorbing particles, enabling polarization-controlled pushing and pulling¹⁵⁰. Shaped hollow beams¹⁵¹ and superpositions of Gaussian and Hermite–Gaussian modes¹⁵² further demonstrate that spatial-mode engineering can improve the confinement and manipulation of absorbing particles. These results show that structured light can tailor not only electromagnetic momentum transfer but also the absorption and temperature distributions that govern heat-mediated particle motion.

4.1.4 Three-Dimensional and Axial Manipulation

The joint intensity–phase strategy described above can be extended from planar trajectories to three-dimensional space by additionally controlling the axial evolution of the structured field. The Grier group designed coordinated three-dimensional intensity distributions and local phase gradients to drive particles along an inclined three-dimensional ring¹⁵³, as shown in Fig. 13(a), and subsequently along a three-dimensional helical trajectory, enabling upward or downward axial transport¹¹², as illustrated in Fig. 13(b).

The Rodrigo group further demonstrated particle transport along a variety of complex curves, including polygons, star-shaped paths, and other closed or open trajectories^{34,154}. Figure 13(c, d) show representative curved optical traps and the corresponding particle motion. These studies also revealed that sharp corners and regions of rapidly changing curvature often produce nonuniform phase gradients in the experimentally generated field. As a result, particles may accelerate, decelerate, or temporarily stall at specific positions¹⁵⁴. For complex trajectories containing sharp turns, the local nonuniformity of the phase gradient must therefore be compensated in addition to designing the geometrical path

itself.

By jointly controlling intensity, phase, and polarization, particle transport can be combined with rotational and orientational control. As illustrated in Fig. 13(e), the intensity gradient confines one or more gold nanorods near a curved trajectory, the phase gradient propels them along the path, and the longitudinal spin angular momentum carried by circularly polarized light¹⁵⁵ controls their rotation and orientation¹⁰³. Orbital motion, self-rotation, and particle orientation can thus be coordinated within a single structured field.

Spatially structured light can also organize optical traps along the propagation direction, thereby providing control over the axial degree of freedom. Owing to their extended nondiffracting range and axially elongated intensity structure, Bessel beams can form stable trapping sites in multiple axial planes. A single Bessel beam has thus been demonstrated to simultaneously trap and manipulate particles at several axial positions⁴⁶. Building on this property, Ruffner and Grier coherently superposed two coaxial Bessel beams to construct a movable periodic array of optical potentials along the propagation direction. As shown in Fig. 14(a), varying the relative phase between the two beams continuously translates the trapping sites along the optical axis. The resulting configuration functions as an “optical conveyor,” enabling active, smooth, bidirectional, and reversible axial transport of particles¹⁵⁶. A limitation of conventional orbital-angular-momentum beams is that their phase handedness normally remains unchanged along the propagation direction. For example, a higher-order Bessel beam has the same sign of topological charge at different axial positions, so trapped particles rotate in the same direction throughout the propagation region. This makes it difficult to independently control particle motion in different axial sections. Real-space phase modulation has been explored as a means of varying orbital angular momentum along the propagation direction^{157–160}, but these schemes generally suffer from degradation of the perfect-vortex property, a limited range of accessible topological charges, or incompatibility with high-numerical-aperture focusing systems. Consequently, they are not readily applicable to precise control of the particle’s axial degree of freedom.

To overcome these limitations, Yuan et al. proposed a longitudinally tunable orbital-angular-momentum vortex beam (LTOVB) designed in the angular-spectrum domain¹⁶¹. Its angular spectrum can be expressed as

$$E_i \left(\sqrt{k^2 - k_z^2}, \varphi \right) = \frac{1}{\text{rect}(k_z/2k)} \int_{-\infty}^{\infty} e^{im(z)\varphi} e^{ik_{z0}(z)z} e^{-ik_z z} dz. \quad (34)$$

Here, $e^{im(z)\varphi}$ represents a vortex phase whose topological charge $m(z)$ varies with the longitudinal position z , thereby enabling axial control of orbital angular momentum. The factor $e^{ik_{z0}(z)z}$ preserves the perfect-vortex characteristics of the field, where $k_{z0}(z)$ is the longitudinal wavevector of the

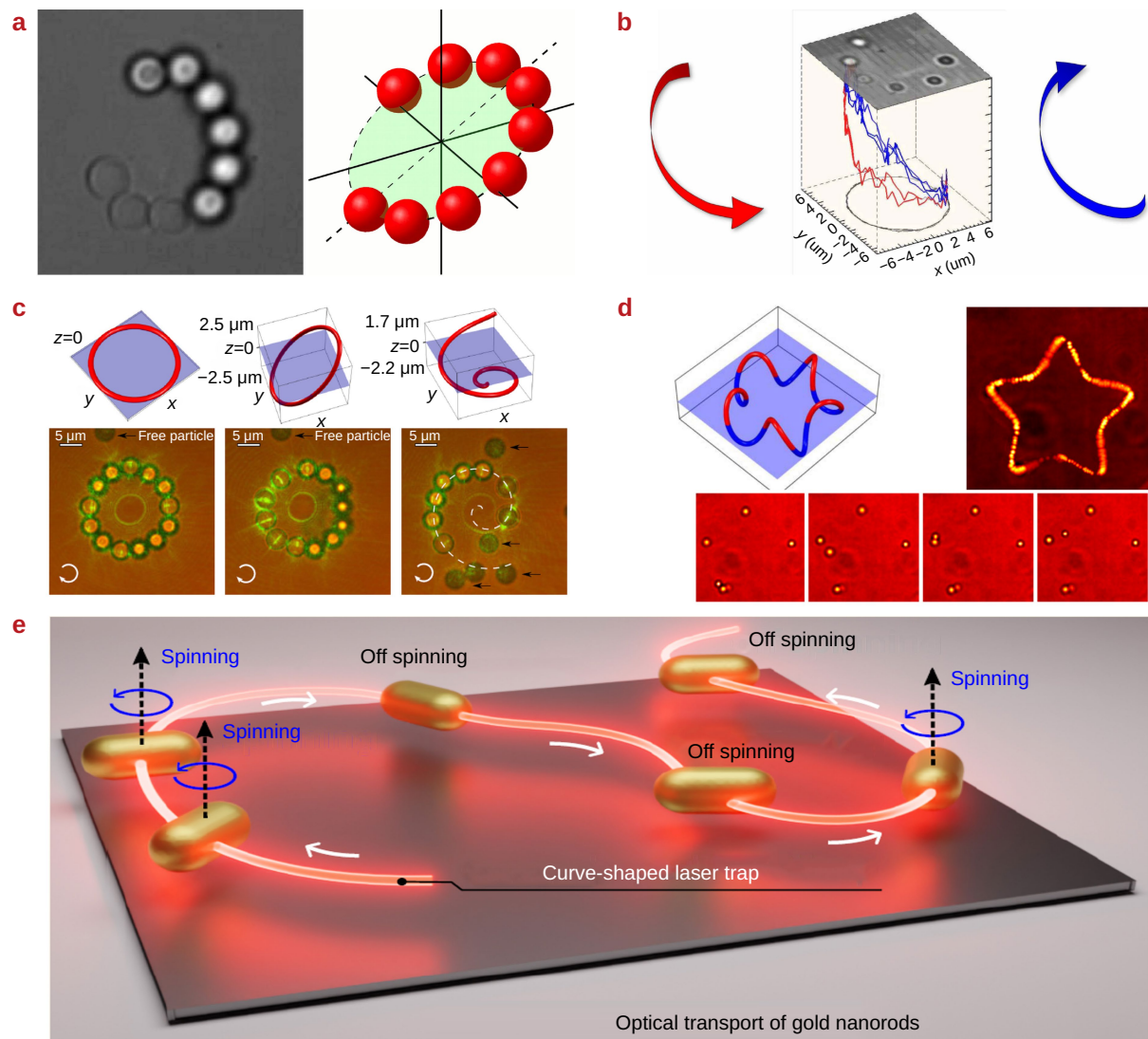


Fig. 13 | Multidimensional manipulation using scalar structured light. (a) Trapping and orbital motion along an inclined three-dimensional ring¹⁵³; (b) bidirectional axial transport along a three-dimensional helical trajectory¹¹²; (c,d) particle transport in various complex curved optical traps^{34,154}; and (e) joint intensity, phase, and circular-polarization control of the orbital motion, self-rotation, and orientation of gold nanorods¹⁰³. Figure reproduced from: (a) ref.¹⁵³, Optica Publishing Group; (b) ref.¹¹², Optica Publishing Group; (c) ref.³⁴, Optica Publishing Group; (d) ref.¹⁵⁴, Optica Publishing Group; (e) ref.¹⁰³, American Chemical Society.

corresponding Bessel mode at position z , satisfying $k_{z0}^2(z) = k^2 - k_{r0}^2(z)$.

As shown in Fig. 14(b), this scheme enables deterministic reversal and switching of the topological charge along the optical axis under tight-focusing conditions, including both two-section and multisection modulation. The direction and speed of particle motion can therefore be controlled independently in different axial regions while the perfect-vortex characteristics of the beam are maintained. This approach introduces an additional degree of freedom for three-dimensional orbital-angular-momentum control and particle manipulation.

4.2 Polarization and Vector Structured-Light Manipulation

Beyond intensity and phase, polarization provides an additional degree of freedom for controlling optical torque and momentum transfer. Recent developments in optical torque have increasingly emphasized two complementary strategies: engineering the incident optical field and tailoring the optical response of the manipulated object¹⁶². We first consider mechanical effects associated with polarization and spin angular momentum, and then extend the discussion to vector structured fields in which the polarization state itself varies spatially.

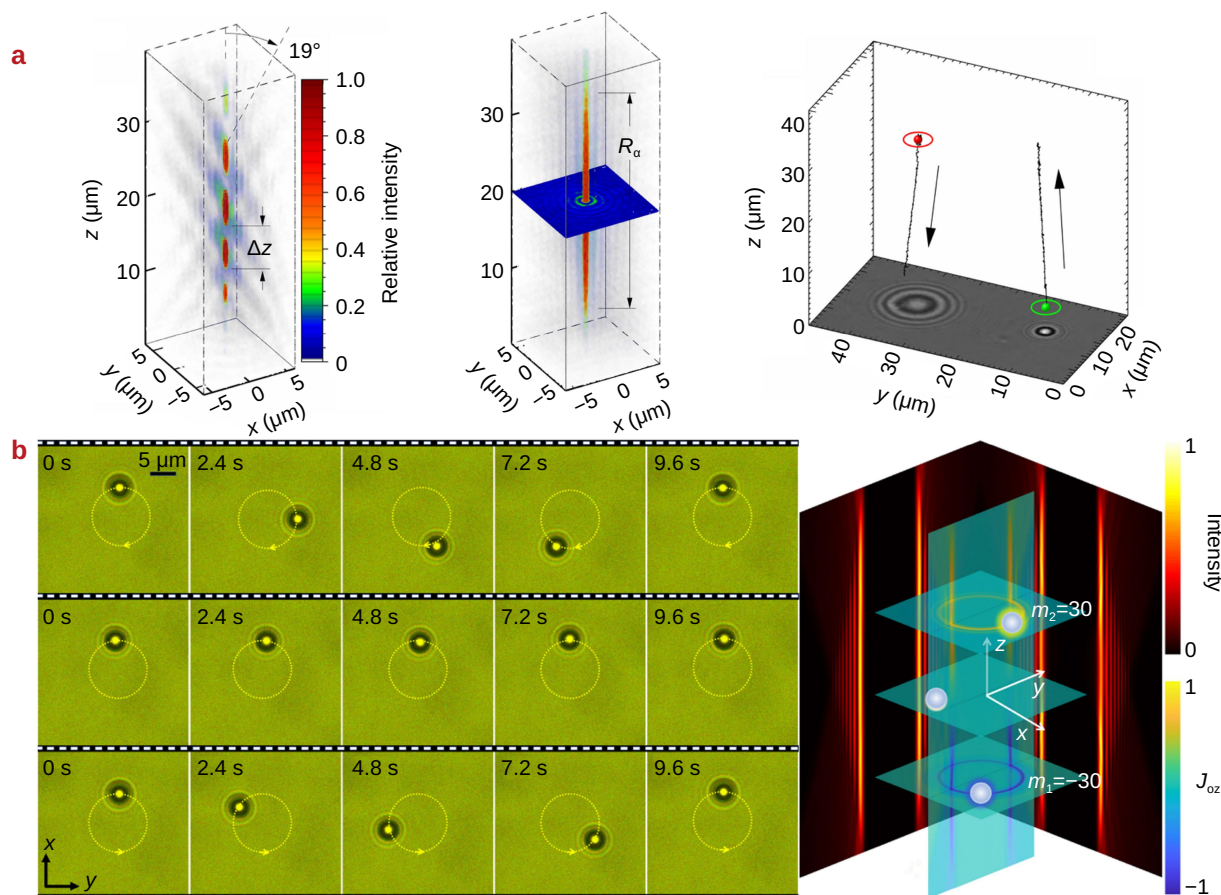


Fig. 14 | Structured-light manipulation along the propagation direction. (a) a movable optical conveyor formed by Bessel beams¹⁵⁶; and (b) control of particle direction and speed in different axial regions using a longitudinally tunable orbital-angular-momentum beam¹⁶¹. Figure reproduced from: (a) ref.¹⁵⁶, American Physical Society; (b) ref.¹⁶¹, Optica Publishing Group.

4.2.1 Polarization- and Spin-Controlled Rotation and Translation

In a paraxial beam, circular polarization carries longitudinal spin angular momentum along the propagation direction and can drive a particle to rotate about an axis parallel to the optical axis^{10,35,163}. Longitudinal spin is associated with rotation of the electric-field vector in a plane perpendicular to the propagation direction. Transverse spin, in contrast, arises when the field vector rotates in a plane containing the propagation direction, resembling the motion of the spokes of a rolling wheel¹⁶⁴. Such transverse spin can drive anisotropic nanoparticles^{165,166} and microparticles¹⁶⁷ to rotate about an axis perpendicular to the optical axis. In one representative configuration, two noncollinear beams with orthogonal circular polarizations interfere such that their longitudinal spin components cancel, while a pronounced transverse spin is reconstructed in the focal region. The resulting transverse spin can drive a micrometer-scale birefringent particle to rotate about a transverse axis¹⁶⁷.

Extending this concept to arbitrary three-dimensional spin control, Wu et al. established an analytical relationship between the local polarization helicity of the incident vector

field and the three-dimensional spin vector of the focused field, enabling spin angular momentum with an arbitrary orientation to be generated in the focal region. As illustrated in Fig. 15, continuous control of the spin-vector direction allows the optical torque acting on an anisotropic particle to be tuned in three dimensions, thereby enabling controlled rotation about an arbitrary axis¹⁶⁸.

In addition to driving self-rotation, spin-related momentum can induce particle orbital motion and transverse translation. As shown in Fig. 16(a), Zhao et al. observed particle orbital motion in a tightly focused circularly polarized field in 2007 and attributed the phenomenon to spin-to-orbital angular-momentum conversion during tight focusing¹⁶⁹. Bekshaev et al. subsequently established a systematic orbital-spin decomposition of optical energy flow and momentum density, providing the theoretical basis for distinguishing the mechanical roles of the two components.

Within this framework, Angelsky et al. performed a key experiment using a two-beam interference configuration with a low numerical aperture, $NA \leq 0.2$. Even under conditions in which spin-to-orbital conversion was negligible,

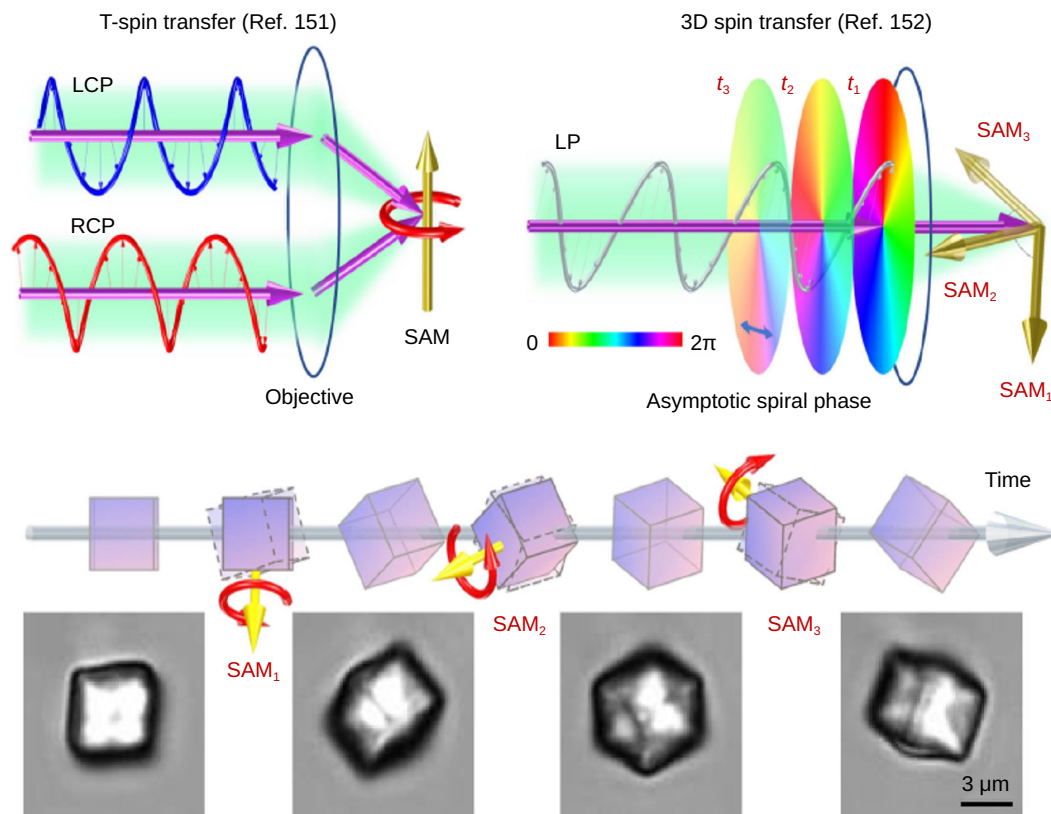


Fig. 15 | Three-dimensional spin-vector control and rotation about an arbitrary axis¹⁶⁸. Figure reproduced from ref.¹⁶⁸, Springer Nature.

particle orbital motion was still observed, demonstrating a measurable mechanical effect associated with the spin momentum density¹⁷⁰. This finding provided an important basis for reinterpreting particle orbital motion in tightly focused circularly polarized beams. In the same year, the group verified the effect in a single Gaussian beam. As shown in Fig. 16(b), reversing the handedness of the incident circular polarization reverses the direction of particle orbital motion around the beam axis¹⁷¹.

Figure 16(c) shows the transverse spin and associated spin-momentum distribution in an evanescent field. Using an evanescent wave as a model system, Bliokh et al. identified the physical origin of transverse spin momentum⁵³, which was subsequently measured experimentally¹⁷². As shown in Fig. 16(d), Huang et al. generated nonzero spin momentum density using an asymmetric radial intensity gradient before and after the focal plane of an axicon and used it to drive particle orbital motion¹⁷³. When partially gold-coated Janus particles were used, the light-avoiding response of the metallic region enabled a dual-orbit mode¹⁷⁴. In addition, Shi et al. utilized transverse spin momentum in a linear structured beam to drive lateral particle translation, as shown in Fig. 16(e)¹⁷⁵. More recently, the mechanical action of spin momentum on particles with general multipolar responses has also been formulated theoretically¹⁷⁶.

4.2.2 Vector Structured-Light Optomechanics

The preceding discussion concerns mechanical effects associated with polarization and spin, which may occur even when the incident polarization is spatially uniform. Vector structured light introduces a further level of control by allowing the polarization orientation, ellipticity, and handedness themselves to vary spatially. These polarization gradients reshape the local momentum and angular-momentum distributions and their spatial curls, giving rise to mechanical responses that are absent or negligible in scalar structured fields.

This section focuses on two classes of momentum effects in vector structured fields. The first involves real momentum and real angular momentum associated with propagating energy flow and momentum exchange. Their spatial distributions are generally related to gradients of the polarization orientation, ellipticity, and handedness. The second concerns imaginary momentum, which has attracted growing interest recently and is determined by the polarization direction, intensity gradient, and their relative angle. For the locally linearly polarized fields considered here, this momentum depends mainly on the intensity gradient and its orientation with respect to the local polarization axis.

For many years, the role of polarization in optical manipulation was considered primarily in relation to anisotropic particles. Linear polarization can orient birefringent

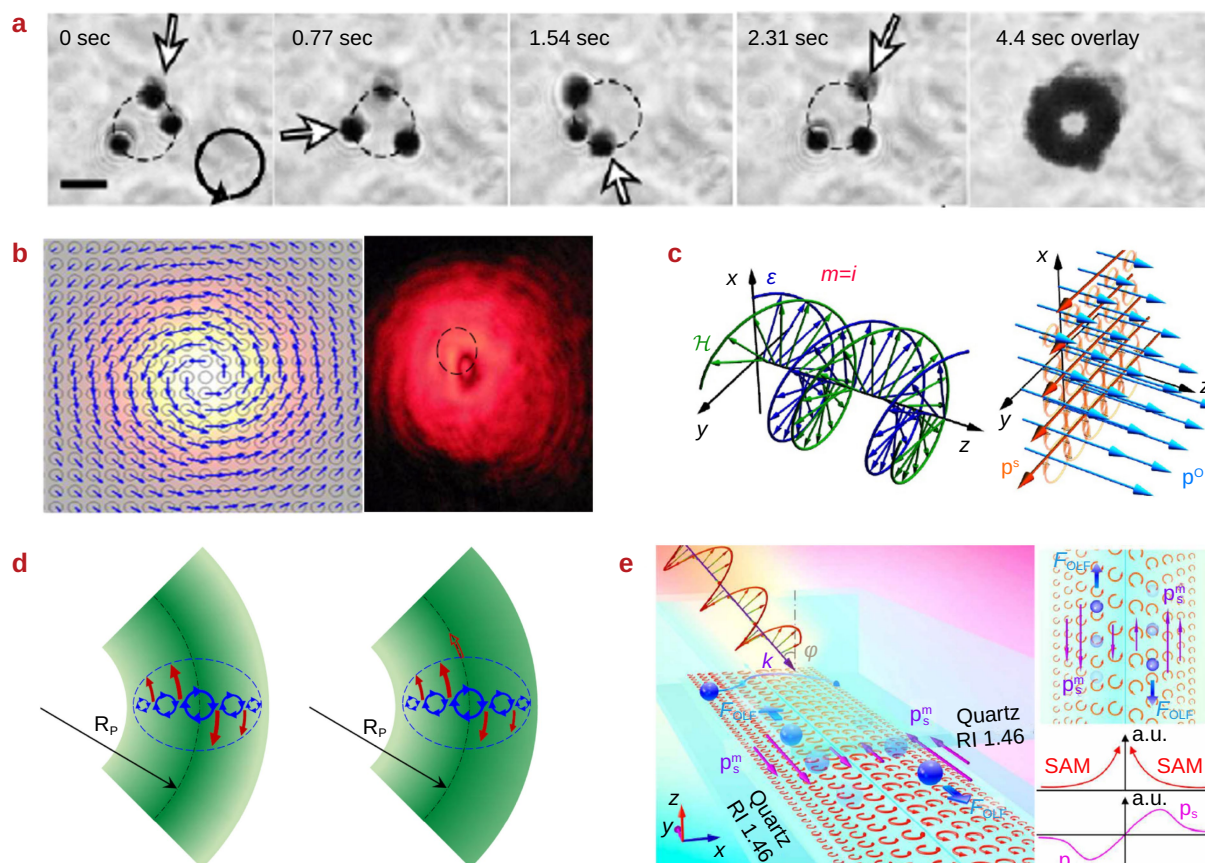


Fig. 16 | Orbital and translational effects associated with spin momentum. (a) Particle orbital motion and spin–orbit coupling in a tightly focused circularly polarized beam¹⁶⁹; (b) spin-momentum-driven orbital motion under weak-focusing conditions¹⁷¹; (c) transverse spin momentum in an evanescent field¹⁷²; (d) spin-dependent orbital motion induced by a radially asymmetric intensity gradient¹⁷³; and (e) transverse translation in a line-shaped structured field¹⁷⁵. Figure reproduced from: (a) ref.¹⁶⁹, American Physical Society; (b) refs.¹⁷¹, Optica Publishing Group; (c) refs.¹⁷², Springer Nature; (d) ref.¹⁷³, Optica Publishing Group; (e) ref.¹⁷⁵, American Association for the Advancement of Science.

objects^{10,177}, whereas circular polarization transfers spin angular momentum and drives their rotation³⁵. The effects of polarization structure on isotropic particles received comparatively little attention.

In 2010, the Cipparrone group first reported a mechanical effect of a vector structured field on isotropic particles¹⁷⁸, demonstrating that a spatial polarization gradient can drive particle motion even when the intensity distribution is nearly unchanged. The vector field shown in Fig. 17(a) was generated through coherent superposition and exhibits a periodic polarization gradient along the horizontal direction, while the polarization remains unchanged along the vertical direction. Fig. 17(b) shows the local polarization structure together with the corresponding force distributions, revealing that the direction of the mechanical action changes periodically with the local polarization orientation. When the entire vector field was translated horizontally, an isotropic particle trapped in the field did not simply follow the field along the translation direction. Instead, it underwent periodic back-and-forth motion in the vertical direction, as shown in Fig. 17(c). This transversely coupled

response experimentally confirmed that a polarization gradient can exert a measurable mechanical effect on an isotropic particle, thereby broadening the role of polarization engineering in optical micromanipulation.

Wang et al. subsequently refined the theoretical formulations of momentum and angular momentum in vector optical fields and experimentally demonstrated that the spatial curl of polarization can generate an effective angular momentum, thereby driving isotropic microparticles along circular trajectories⁶⁶. As shown in Fig. 18(a), particles remain stationary when the vector field consists of a locally linearly polarized structure. By contrast, when the polarization ellipticity varies periodically along the radial direction, the particles begin to orbit. This result challenged the conventional view that polarization primarily acts on anisotropic particles and demonstrated that radial gradients of the polarization state in vector structured fields can reshape the momentum and angular momentum distributions of the optical field. Pan et al. subsequently demonstrated experimentally that angular momentum induced by the azimuthal winding gradient of the polarization orienta-

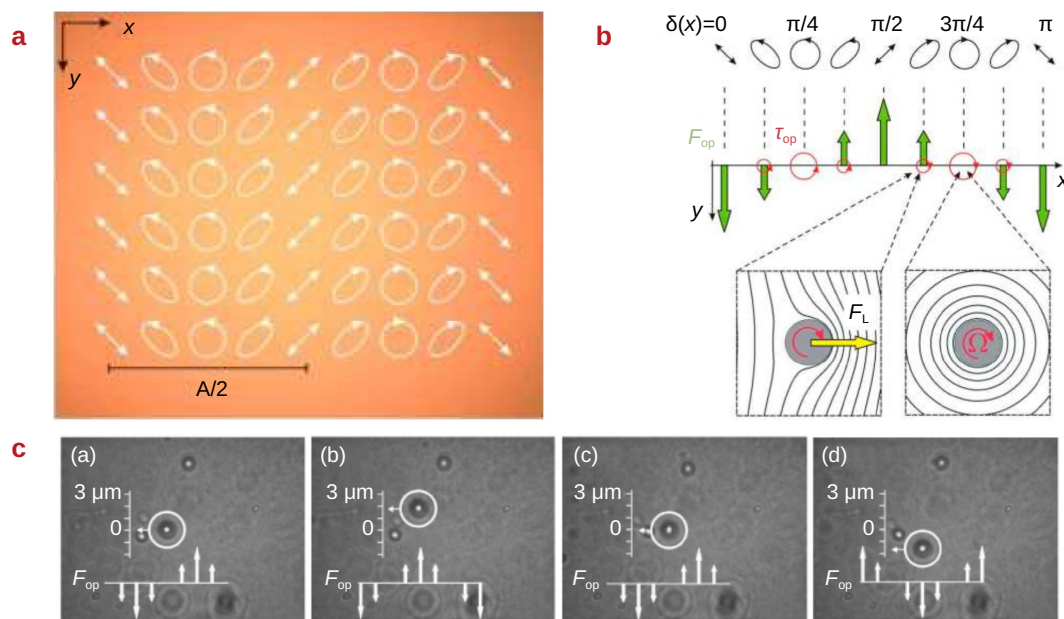


Fig. 17 | Particle motion induced by a polarization gradient¹⁷⁸. (a) A vector field with polarization varying periodically along the horizontal direction; (b) the local polarization structure and corresponding force analysis; and (c) periodic transverse motion of an isotropic particle when the field is translated. Figure reproduced from ref.¹⁷⁸, Optica Publishing Group.

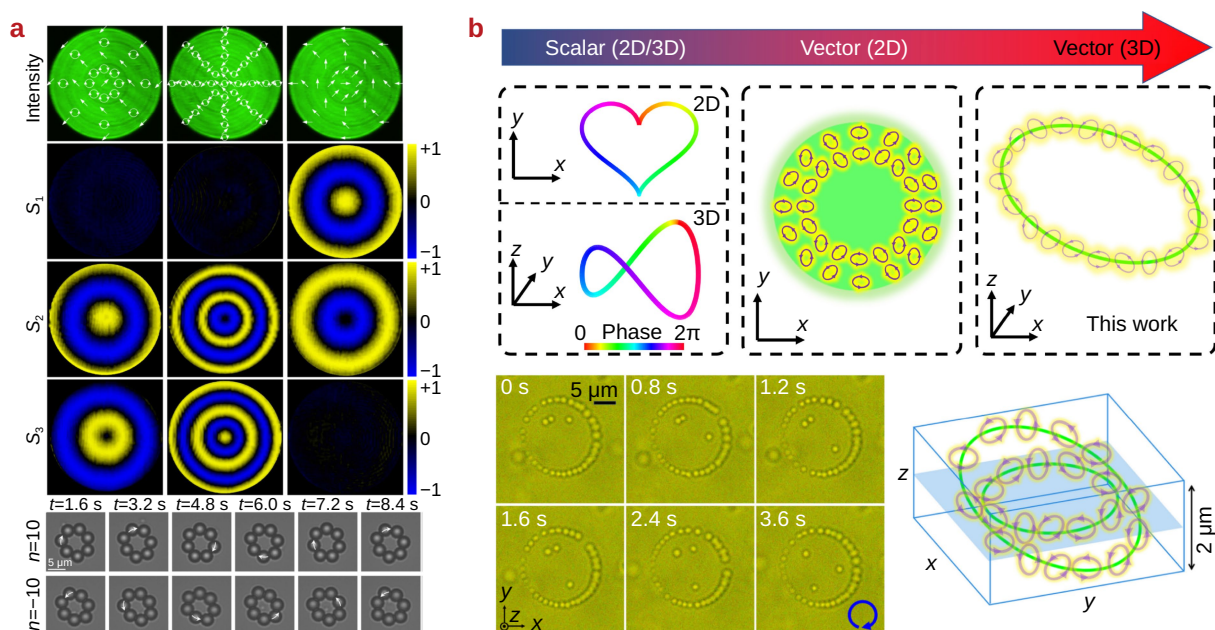


Fig. 18 | Angular momentum and orbital motion in vector structured fields. (a) Orbital motion of isotropic particles driven by a radial gradient in polarization ellipticity⁶⁶. (b) Three-dimensional transport of particles in a vectorial curved optical field⁶⁵. Figure reproduced from: (a) ref.⁶⁶, American Physical Society; (b) ref.⁶⁵, Optica Publishing Group.

tion can drive particle motion¹⁷⁹. Yuan et al. further extended this mechanism into three dimensions⁶⁵. As shown in Fig. 18(b), continuously engineering the polarization orientation and its winding gradient along a three-dimensional curve generates an angular momentum flow distributed along the trajectory, thereby transporting particles along arbitrary three-dimensional paths.

The studies described above primarily concern the real part of the complex momentum density. Real momentum generally corresponds to propagating energy flow and produces translational or orbital motion through momentum exchange between the optical field and the particle. The imaginary momentum, by contrast, is more closely associated with local reactive fields and electromagnetic energy

storage and was long regarded as having little direct mechanical significance. Recent theoretical studies have shown that imaginary momentum can contribute to measurable time-averaged forces through suitable electric–magnetic or multipolar scattering channels^{55,67}.

Zhou et al. provided the first experimental observation of particle motion driven by the imaginary momentum of a vector structured field¹⁸⁰. As shown in Fig. 19(a), they constructed a cylindrical-vector double-ring field whose real momentum was directed predominantly along the propagation axis and contained essentially no azimuthal component. The imaginary momentum, however, exhibited a pronounced azimuthal component. Its magnitude depended on the angle between the local linear-polarization direction and the intensity gradient and reached a maximum when this angle was approximately 45°.

Within the double-ring field, the imaginary momentum in the inner and outer rings was distributed symmetrically but circulated in opposite directions. Particles trapped in the two rings therefore orbited oppositely: one clockwise and the other counterclockwise. The experimental results agreed with theoretical predictions and demonstrated that imaginary momentum can produce a distinct and measurable mechanical contribution.

Imaginary-momentum effects were subsequently verified in flat-top vector beams⁶⁸. As shown in Fig. 19(b), changing the angle between the polarization direction and the intensity gradient reverses the direction of the imaginary momentum, thereby driving directional particle translation and controlling the rotational direction of a nanorod. Imaginary-momentum-induced motion in Gaussian beams¹⁸¹, as well as fixed-point trapping and transport assisted by imaginary momentum¹⁸², have also been reported, further supporting the generality and control potential of this mechanism.

Chirality provides another particle-response degree of

freedom for structured-light manipulation. Chiral optical fields and spatially tailored polarization can generate enantioselective or lateral optical forces, enabling the separation and sorting of particles with opposite handedness^{57–59}. Recent advances have further extended chirality-sensitive optical manipulation toward integrated sorting and detection schemes, highlighting the potential of structured optical fields for selective manipulation and characterization of chiral micro- and nanoscale objects¹⁸³.

4.3 Engineered Particles and Light-Driven Micro-/Nanorobots

The approaches discussed above primarily tailor the incident optical field while treating the particle response as given. A complementary strategy is to engineer the optical properties or internal structure of the manipulated object itself, thereby reshaping its momentum exchange with light. Recent experiments have shown, for example, that an inhomogeneous chirality distribution can generate a reversible lateral optical force and induce self-swerving motion of chiral particles even under fixed linearly polarized illumination¹⁸⁴. This result illustrates that complex mechanical responses can arise not only from structuring the incident field but also from spatially engineering the light–matter response of the object.

Using a centimeter-scale element with spatially structured birefringence, Magallanes and Brasselet directly observed a transverse displacement that reversed with the incident optical spin, together with a left-handed optical torque opposite to the incident spin angular momentum¹⁸⁵. This macroscopic experiment demonstrated that engineering the internal polarization response and scattering structure of an object can reverse or redirect the resulting optical recoil force and torque.

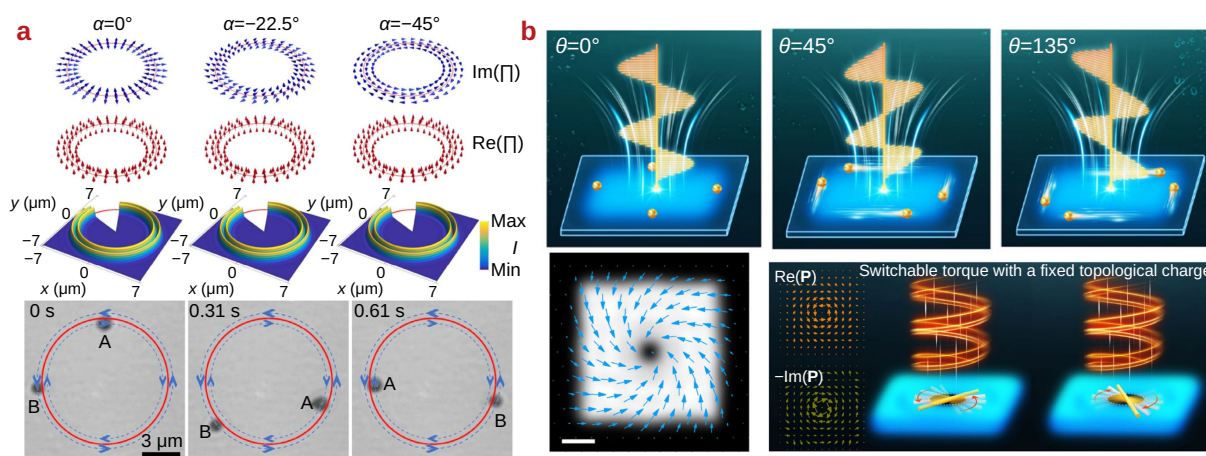


Fig. 19 | Imaginary momentum and its mechanical effects. (a) Imaginary-momentum distribution and experimental observation in a cylindrical-vector double-ring field¹⁸⁰; and (b) particle translation and nanorod rotation driven by imaginary momentum in a flat-top vector beam⁶⁸. Figure reproduced from: (a) ref.¹⁸⁰, National Academy of Sciences; (b) ref.⁶⁸, Springer Nature.

At the microscale, the same concept has developed into light-driven vehicles capable of autonomous propulsion and steering. Andr  n et al. introduced the concept of optical “metavehicles”¹⁸⁶, in which a dielectric metasurface with directional-scattering functionality is embedded within a microparticle. Under plane-wave illumination, asymmetric scattering generates recoil momentum that propels the vehicle, while changes in the incident polarization modify the response of different metasurface regions and thereby control its direction and trajectory. Wu et al. subsequently developed light-driven “microdrones”¹⁸⁷ by integrating four independently addressable chiral plasmonic nanomotors into a microscale polymer carrier. The left- and right-circularly polarized components at two operating wavelengths provide four independent control channels. By independently tuning the recoil forces generated by the nanomotors,

the two translational degrees of freedom and one rotational degree of freedom in the plane can be controlled separately, enabling flexible translation, rotation, and navigation along complex trajectories.

Qin et al. further integrated a cross-antenna plasmonic nanotweezer with localized near-field enhancement into the microdrone platform¹⁸⁸. Figure 20(a) shows the resulting microrobot, which combines independently addressable plasmonic motors for propulsion and steering with an integrated nanotweezer for nanoparticle manipulation. Simultaneous illumination at 830 and 980 nm enables translational motion while maintaining the trapping of a single 70 nm fluorescent nanodiamond, allowing the particle to be transported and reversibly released. The same group subsequently miniaturized the system to the submicrometer scale and developed a light-driven nanorobotic cleaner¹⁸⁹.

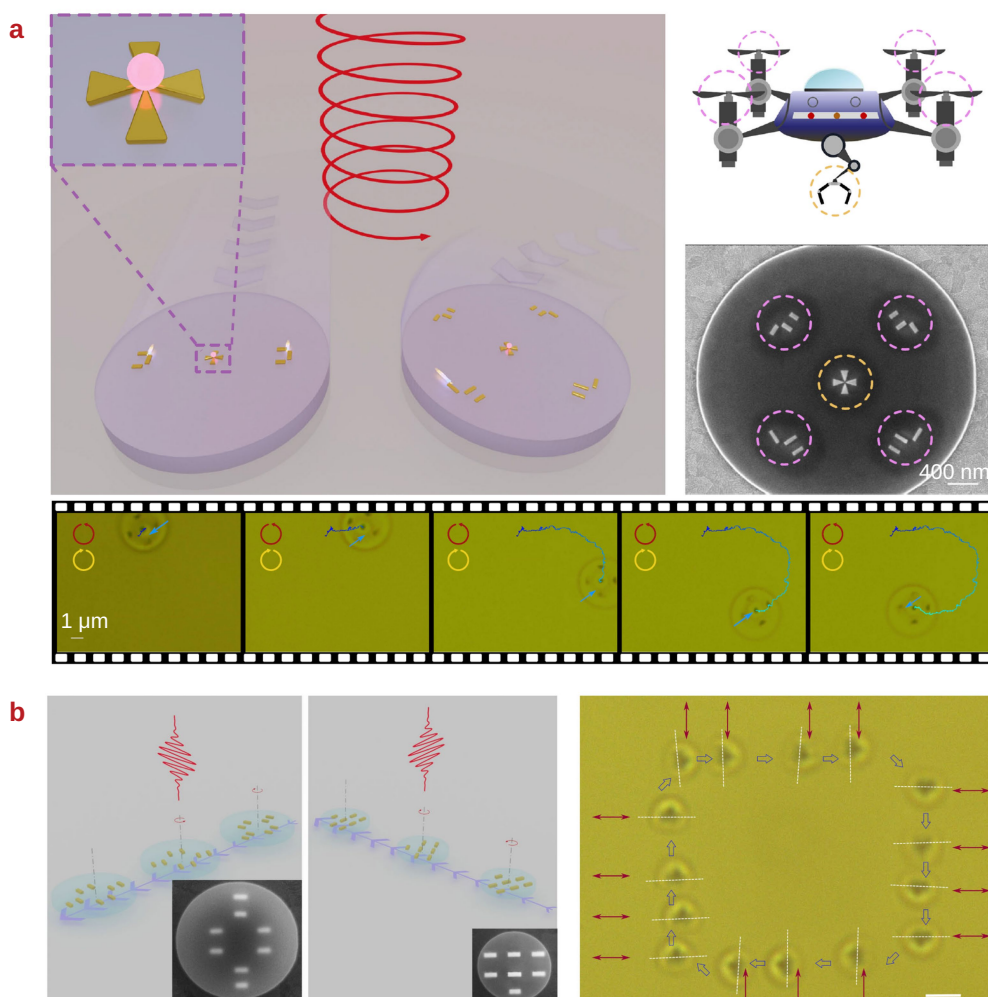


Fig. 20 | Light-driven plasmonic micro- and nanorobots. (a) A plasmonic microrobot integrating independently addressable nanomotors with a cross-antenna nanotweezer; the lower sequence shows the trapping, transport, and reversible release of a 70-nm fluorescent nanodiamond under 830- and 980-nm illumination¹⁸⁸. (b) Micro- and nanorobot designs and their polarization-controlled actuation. Switching between orthogonal linear polarizations produces rectangular motion, while brief circularly polarized illumination induces turns¹⁸⁹. Figure reproduced from: (a) ref.¹⁸⁸, Springer Nature; (b) ref.¹⁸⁹, Springer Nature.

Figure 20(a) illustrates the micro- and nanorobot designs and their polarization-controlled actuation. The microrobot employs separate structures for propulsion and orientation locking, whereas the nanorobot uses a compact single-motor configuration that simultaneously generates optical thrust and a restoring torque. Under unfocused 980 nm illumination, both designs remain aligned with the incident linear polarization and exhibit controlled translational motion. By switching between two orthogonal linear polarizations and briefly applying circularly polarized illumination at each corner, the nanorobot can follow a rectangular trajectory with deterministic turns. It can also capture, assemble, transport, and release bacteria through opto-thermophoretic interactions, enabling it to function as a light-driven robotic cleaner.

Beyond plasmonic and metasurface-based vehicles, light-driven microrobotic systems have also been developed using dielectric, birefringent, and biohybrid structures. Dielectric microrobots fabricated by microstructuring can be actuated by distributed optical forces for multidimensional pose control and indirect manipulation of micro-objects¹⁹⁰. Form-birefringent microgears provide another route for converting polarization into mechanical rotation¹⁹¹, while optically driven birefringent microspheres can generate localized microfluidic shear and have been used to steer nerve-fibre growth¹⁹². Biohybrid systems further exploit the intrinsic phototaxis of microorganisms, as demonstrated by light-controlled *Volvox* microrobots for directional cargo transport in microchannels¹⁹³. These examples broaden light-driven microrobotics from engineered optical recoil to mechanisms based on distributed optical forces, birefringent torque, and biological photoresponses.

5 Conclusions and Outlook

This review has examined the physical foundations, field-generation methods, and representative applications of spatially structured light in optical micromanipulation. By tailoring the spatial distributions of amplitude, phase, and polarization, the intensity gradient, local energy flow, and linear and angular momentum distributions within the interaction region can be reshaped, enabling the on-demand control of optical forces, optical torques, and trapping potentials. Starting from the orbital-spin decomposition of optical momentum, we have discussed the real and imaginary momentum components and their relationships with the electromagnetic responses of particles, and traced the development of scalar and vector structured optical fields from two-dimensional complex-amplitude modulation to three-dimensional fully vectorial field synthesis. Overall, spatially structured light is driving optical micromanipulation from single-point, static, and low-dimensional control toward multiparticle, programmable, and multidimensional manipulation. Polarization gradients, spin-related momentum, and imaginary momentum have further enriched the

mechanisms available for driving particle motion, while the integration of metasurfaces, plasmonic nanomotors, and near-field nanotweezers into manipulated objects has advanced light-driven systems from passively controlled particles toward functional micro- and nanorobots capable of propulsion, navigation, orientation control, and material handling.

From the perspective of fundamental physics, establishing a quantitative correspondence between complex three-dimensional structured optical fields and particle mechanical responses remains an important challenge. As optical-field engineering advances from two-dimensional scalar distributions to three-dimensional fully vectorial structures, the local amplitude, phase, and polarization all vary throughout space, while high-numerical-aperture focusing and particle scattering can further reshape the momentum and angular momentum distributions in the focal region. High-precision, real-time characterization of complex three-dimensional vector fields, together with quantitative mapping between experimentally measured local field structures and the resulting optical forces and torques, will therefore be essential for achieving precise manipulation with a large number of degrees of freedom. Meanwhile, structured-light manipulation is progressing from static spatial shaping toward joint spatiotemporal control and ultrafast optics^{194–195}. Spatiotemporal optical vortices carrying transverse orbital angular momentum offer new degrees of freedom for non-coaxial transport, transverse rotation, and transient optical-torque control¹⁹⁶, although how such ultrafast spatiotemporal structures can be efficiently converted into measurable and controllable microscopic mechanical responses remains to be explored. Another fundamental challenge arises when moving from single-particle manipulation to dense multiparticle and complex collective systems. Multiple scattering, optical binding, and hydrodynamic coupling among particles make the relationship between local optical forces and the resulting collective dynamics substantially more complicated. Developing a unified framework that extends single-particle optical-force descriptions to many-body interactions and collective dynamics will therefore be important for large-scale self-organization and cooperative manipulation with structured light.

From the perspective of device and control technologies, high-speed, precise, and high-dimensional dynamic manipulation remains jointly constrained by field-computation speed, modulator refresh rate, and feedback latency. Current spatial light modulators offer considerable flexibility for generating complex structured fields, but significant trade-offs remain among spatial resolution, pixel count, field of view, modulation efficiency, and refresh rate. Optically addressed metasurface spatial light modulators have already demonstrated dynamic complex-amplitude holography, providing a promising route toward compact and high-speed optical-field reconfiguration¹⁹⁷. Further integration of dynamic metasurfaces, integrated photonics, and highly effi-

cient structured-light generators may reduce system complexity while increasing modulation bandwidth.

From the perspective of interdisciplinary applications, structured-light micromanipulation should continue to expand across different wave systems, manipulation environments, and target objects, while addressing the physical constraints specific to each platform. The concepts of structured-wave manipulation have already extended beyond optics to acoustic fields and water waves. Holographic acoustic fields have enabled object levitation, translation, and rotation¹⁹⁸, while topological water waves have been used to trap, orbit, and spin floating objects¹⁹⁹. Manipulation environments are also expanding from liquids and gases toward solid-state and phase-transition systems. For example, laser-induced local melting and thermocapillary convection can drive particle transport and device assembly within solid-cladding optical fibers²⁰⁰, whereas ice-phase optothermal tweezers exploit localized melting at ice–water interfaces to achieve programmable transport and assembly of micro- and nanoscale objects within ice matrices²⁰¹. These developments suggest that the concepts of spatial energy-flow and momentum engineering established in structured light may evolve toward a more general framework for structured-wave manipulation. Nevertheless, dissipation, interfacial effects, and nonlinear responses in different media will require system-specific theoretical descriptions and control strategies.

The range of manipulable objects is also expected to expand from classical particles toward atoms, quantum systems, and increasingly complex biological objects. Reconfigurable atomic arrays provide important platforms for quantum simulation and quantum computing²⁰², while the integration of optical tweezers with microfluidics and the development of optical tweeze-sectioning microscopy offer new capabilities for cell sorting, precision manipulation, three-dimensional imaging, and array organization of suspended cells^{203–204}. In contrast to inorganic particles, however, long-term optical manipulation of living cells and biomolecules remains constrained by local heating, photochemical damage, limited penetration depth, and the complexity of biological environments^{205–206}. Reducing optical power and local energy deposition, improving trapping efficiency, and combining longer-wavelength illumination with low-damage structured fields and real-time feedback will therefore be important for improving biocompatibility and extending optical manipulation toward *in vivo* environments²⁰⁷. As these challenges in fundamental physics, device technology, and interdisciplinary applications are progressively addressed, structured-light micromanipulation may evolve from precise control of individual particles toward high-dimensional, adaptive, and intelligent manipulation platforms for complex many-body systems, functional micro- and nanorobots, and biological and quantum systems.

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Author contributions

Z. Yuan and J. Ding wrote and revised the manuscript. J.Ding supervised the project.

Competing interests

The authors declare no competing financial interests.



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