Dynamic polarization-dependent optical-vortex-controlling via a fiber with acousto-optic interaction

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Abstract. We report a new dynamic polarization-dependent mode conversion in a circular optical fiber endowed with the acousto-optic interaction, in which the sign of the topological charge of the generated fiber optical vortex are governed by the linear polarization direction of the incident beam. The described effect can be useful in developing new acousto-optic devices such, for example, as polarization-controlled optical vortex intensity modulators. Yet, this type of all-fiber wavelength-tunable optical vortex generation and controlling should be especially useful in such vortex-based applications as micromechanics, classical and quantum information encryption, and simulation of quantum computing.

1. Introduction

The possibility of interaction between light and acoustic waves - the so-called acousto-optic interaction (AOI) - was first reported in a seminal paper by Brillouin [1] in 1922. From the classical point of view, the AOI comes from the light diffraction on a moving refractive index grating formed by an acoustic perturbation in a medium. Quantum mechanically, the photon-phonon scattering is shown to be the underlying physics. From the laws of energy $\omega' = \omega + \Omega$ and momentum $\mathbf{k}' = \mathbf{k} + \mathbf{K}$ conservation, where $(\omega, \mathbf{k})$, $(\omega', \mathbf{k}')$ and $(\Omega, \mathbf{K})$ are the frequency and wave-vector of the incident photon, the diffracted photon and the phonon, respectively, it follows that the AOI can be efficiently used for controlling of the light propagation. Indeed, frequency shifters, wavelength tunable filters, deflectors and intensity modulators have been implemented both in the bulk [2] and fiber configurations[3, 4, 5, 6, 7, 8, 9]. The last case is special because fibers guide both acoustic [10, 11] and optical beams [12] that provides unique conditions for a highly efficient AOI. In addition, a fiber scheme offers valuable advantages of the all-fiber fast dynamic light manipulation and low insertion loss.

It should be noted that the flexural acoustic wave (FAW) is primarily used in fiber acousto-optics because it produces an axially-asymetric perturbation, which enables the well-known coupling of the input fundamental mode LP_0 to the higher-order LP_1 modes [3, 8]. As a recent important development, the researches unveiled the possibility of generating and controlling topologically-charged fiber optical vortices (OVs) [13], which carry orbital angular momentum (OAM) [14], via the AOI induced by the FAW of both the lowest-order [15, 16, 17] and the higher-order [18] in the form of acoustic vortices [19].

Here we aim at demonstrating that a conventional circular fiber endowed with the lowest-order FAW can be used for dynamically-controlled wavelength-tunable highly efficient flipping of OAM of an incident OV with a unity topological charge.

2. Resonance fiber modes

The permittivity of a fiber under consideration is given by:

$$\hat{\varepsilon}(r,\varphi,z,t)=\varepsilon_0(r)+2\Delta\varepsilon_{co}u_0 f_r\cos\varphi\ \cos(Kz-\Omega t)+\varepsilon_{\rm co}^2 pKu_0\begin{pmatrix}0&0&1\\0&0&0\\1&0&0\end{pmatrix}\sin(Kz-\Omega t)\,. \quad (1)$$

Here the unperturbed fiber permittivity $\varepsilon_0(r)=\varepsilon_{\rm co}[1-2\Delta f(r)]$, where $\Delta=(\varepsilon_{\rm co}-\varepsilon_{\rm cl})/2\varepsilon_{\rm co}$ is the normalized index difference, $\varepsilon_{\rm co}$ and $\varepsilon_{\rm cl}$ are the core and cladding values of the permittivity, respectively, and $f(r)$ is the fiber's profile function. In the second term $f_r=df/dr$ and u_0, K, and Ω are the amplitude, wavevector and frequency of the FAW, respectively. In the last term for silica at the wavelength $\lambda=0,\,63\ \mu m$ the constant of photoelasticity is $p=-0.075$. The cylindrical coordinates $(r,\,\varphi,\,z)$ are implied.The following expressions of the modes of a fiber model in the problem are known to be:

$$\begin{aligned}
|\Psi_1^{(\sigma)}\rangle &= \Big[\sin\theta|{\rm LP}_0^\sigma\rangle+\cos\theta|{\rm LP}_1^{ev,\,\sigma}\rangle e^{i(\Omega t-Kz)}\Big]e^{i(\beta_1^\sigma z-\omega t)}\,,\\
|\Psi_2^{(\sigma)}\rangle &= \Big[\cos\theta|{\rm LP}_0^\sigma\rangle-\sin\theta|{\rm LP}_1^{ev,\,\sigma}\rangle e^{i(\Omega t-Kz)}\Big]e^{i(\beta_2^\sigma z-\omega t)}\,,\\
|\Psi_3^{(\sigma)}\rangle &= \Big[\sin\theta|{\rm LP}_1^{ev,\,\sigma}\rangle-\cos\theta|{\rm LP}_0^\sigma\rangle e^{i(Kz-\Omega t)}\Big]e^{i(\beta_3^\sigma z-\omega t)}\,,\\
|\Psi_4^{(\sigma)}\rangle &= \Big[\cos\theta|{\rm LP}_1^{ev,\,\sigma}\rangle+\sin\theta|{\rm LP}_0^\sigma\rangle e^{i(Kz-\Omega t)}\Big]e^{i(\beta_4^\sigma z-\omega t)}\,,\\
|\Psi_5^{(\sigma)}\rangle &= |{\rm LP}_1^{od,\,\sigma}\rangle e^{i(\tilde\beta_1 z-\omega t)}\,. \qquad (2)
\end{aligned}$$

Here in the basis of linear polarizations $|\Psi_k^{(\sigma)}\rangle=(E_x,E_y)^{\rm T}$ the standard even and odd LP modes read as $|{\rm LP}_\ell^{ev,\,x}\rangle=\sqrt{2}F_\ell(r)(\cos\ell\varphi,0)^{\rm T}$, $|{\rm LP}_\ell^{ev,\,y}\rangle=\sqrt{2}F_\ell(r)(0,\cos\ell\varphi)^{\rm T}$, $|{\rm LP}_\ell^{od,\,x}\rangle=\sqrt{2}F_\ell(r)(\sin\ell\varphi,0)^{\rm T}$, $|{\rm LP}_\ell^{od,\,y}\rangle=\sqrt{2}F_\ell(r)(0,\cos\ell\varphi)^{\rm T}$, where $\sigma=x,y$ specifies the direction of linear polarization, $F_\ell(r)$ is the well-known radial function [12], the radial number is omitted and T stands for the transposition. The energy distribution within the hybrid modes is governed by the parameter $0<\theta\le\pi/4$ defined as $\cos 2\theta=(\epsilon/\sqrt{\epsilon^2+Q^2})$. Here $\epsilon=K-\bar K$ and the resonance value of the acoustic wave vector $\bar K=\tilde\beta_0-\tilde\beta_1$ is defined through the well-known scalar propagation constants $\tilde\beta_\ell$ [12] of the ${\rm LP}_{0,1}^\sigma$ modes, respectively. The parameter Q, which characterizes the coupling strength between the unperturbed modes is:

$$Q=Q_g+Q_p\,,Q_g=\sqrt{\frac{\varepsilon_{\rm co}}{2r_0^2N_0N_1}}k\Delta u_0\,,Q_p=-\sqrt{\frac{\varepsilon_{\rm co}^2}{8r_0^2N_0N_1}}pKu_0\int_0^\infty R\frac{dF_0}{dR}F_1dR\,, \qquad (3)$$

where $k=2\pi/\lambda$ and the mode normalization $N_\ell=\int_0^\infty RF_\ell^2dR$. Here we imply the conventional step-index fibers with profile function $f(r)=\Theta(r/r_0-1)$, where Θ is the unit step function. The propagation constants of modes (2) are found to be:

$$\begin{aligned}
\beta_1^{x,y} &= \tilde\beta_0+(1/2)(\epsilon+\sqrt{\epsilon^2+Q^2})\pm(Q_p/2)\sin 2\theta\,,\\
\beta_2^{x,y} &= \tilde\beta_0+(1/2)(\epsilon-\sqrt{\epsilon^2+Q^2})\mp(Q_p/2)\sin 2\theta\,,\\
\beta_3^{x,y} &= \tilde\beta_1+(1/2)(-\epsilon-\sqrt{\epsilon^2+Q^2})\mp(Q_p/2)\sin 2\theta\,,\\
\beta_4^{x,y} &= \tilde\beta_1+(1/2)(-\epsilon+\sqrt{\epsilon^2+Q^2})\pm(Q_p/2)\sin 2\theta\,. \qquad (4)
\end{aligned}$$

Here the upper sign at the last terms corresponds to the x-polarized modes in Eq. (2).

Now, let the OV linearly polarized along x or y-axis $|\mathrm{LV}_\ell\rangle \stackrel{\sigma}{=} |\mathrm{LP}_1^{ev,\sigma}\rangle + i|\mathrm{LP}_1^{od,\sigma}\rangle$ at frequency ω be propagating in the fiber with the above described AOI:

$$|\Psi_{in}\rangle = |\mathrm{LV}_\ell^\sigma\rangle e^{-i\omega t}\,. \tag{5}$$

It will excite in the fiber the superposition of eigenstates (2) with the same polarization σ: $|\Psi(z)\rangle = \sum_k a_k |\Psi_k^{(\sigma)}(z)\rangle$. Invoking the simplified boundary conditions: $|\Psi_{in}\rangle = \sum_k a_k |\Psi_k^{(\sigma)}(z=0)\rangle$, one can determine the decomposition coefficients a_k and the field in the AOI region, which within a phase factor, can be brought to the following form:

$$|\Psi(z)\rangle = \Big[c_0^\sigma(z)|\mathrm{LP}_0^\sigma\rangle e^{i(Kz-\Omega t)} + c_1^\sigma(z)|\mathrm{LV}_\ell^\sigma\rangle + c_2^\sigma(z)|\mathrm{LV}_{-\ell}^\sigma\rangle\Big] e^{-i\omega t}, \tag{6}$$

where the coefficients $c_\ell(z)$ are given by:

$$\begin{aligned} &c_0^\sigma(z) = (i/2)\sin 2\theta \sin(\eta_\sigma z)\,, \\ &c_\ell^\sigma(z) = (1/2)\{[\cos(\eta_\sigma z) + i\cos 2\theta \sin(\eta_\sigma z)]e^{-0.5i\epsilon z} + 1\}\,, \\ &c_{-\ell}^\sigma(z) = (1/2)\{[\cos(\eta_\sigma z) + i\cos 2\theta \sin(\eta_\sigma z)]e^{-0.5i\epsilon z} - 1\}\,, \end{aligned} \tag{7}$$

with $\eta_{x,y} = 0.5(\sqrt{\epsilon^2 + Q^2} \pm Q_p \sin 2\theta)$. Note that $|c_0^\sigma|^2 + |c_1^\sigma|^2 + |c_2^\sigma|^2 = 1$. Eq. (6) shows that the field in the fiber is composed of the incident OV mode $|\mathrm{LV}_\ell^\sigma\rangle$ and a generated OV $|\mathrm{LV}_{-\ell}^\sigma\rangle$ with opposite sign of topological charge as well as the frequency upshifted fundamental mode $|\mathrm{LP}_0^\sigma\rangle$. The energy of the partial beams is determined as:

$$W_\ell^\sigma = |c_\ell^\sigma|^2\,. \tag{8}$$

When (i) the resonance regime $\epsilon = 0$ is implemented and (ii) the fiber has the optimal length $z = L_m^\sigma$ (see figure 1)s, where

$$L_m^x = \frac{2(2m+1)\pi}{Q_g + 2Q_p}\,,\ L_m^y = \frac{2(2m+1)\pi}{Q_g}\,, \tag{9}$$

$W_2^\sigma(0, L_m^\sigma) = 1$, all the incident energy becomes stored in the generated OV mode with the topological charge opposite to the incident one (see figure 1).

$$|\mathrm{LV}_\ell^\sigma\rangle\ e^{-i\omega t} \longrightarrow |\mathrm{LV}_{-\ell}^\sigma\rangle\ e^{-i\omega t}\,. \tag{10}$$

It is important to note that the corresponding conversion length is polarization-dependent, as is seen from Eq. (9). Since $Q_g/Q_p \sim 2k\Delta/n_{\mathrm{co}}pK$ and $\cancel{K} \sim \overline{2\Delta}/r_0$ [12], one get $Q_g/Q_p \sim 100\sqrt{\Delta}(r_0/\lambda) \gg 1$ for all reasonable parameters of weakly-guiding fibers, so that the relative difference in conversion lengths is small. For example, at the waveguide parameter $V = 4.16$, the normalized index difference $\Delta = 0.001$, the core radius $r_0 = 6.3\ \mu m$ and the acoustic power $P = 50\,\mathrm{mW}$, one has: $L_0^x = 4.26$ cm, $L_0^y = 3.56$ cm, hence $L_0^x - L_0^y = 7$ mm.

Nevertheless, such a subtle difference between conversion lengths may be relevant at the corresponding fiber lengths. Indeed, as can be easily shown, at the fiber's length $z = L_k$ given by:

$$L_k = \frac{4\pi}{Q_g}\left\lceil (k + 1/2)\frac{Q_g}{2|Q_p|}\right\rceil\,, \tag{11}$$

where $\lceil x \rceil$ is the ceil function and $k = 0, 1, \ldots$, one gets $W_{-\ell}^x(0, L_k) \approx W_\ell^y(0, L_k) \approx 1$ (see figure 1), which entails:

$$\begin{aligned} &|\mathrm{LV}_\ell^x\rangle\ e^{-i\omega t} \longrightarrow |\mathrm{LV}_{-\ell}^x\rangle\ e^{-i\omega t}\,, \\ &|\mathrm{LV}_\ell^y\rangle\ e^{-i\omega t} \longrightarrow |\mathrm{LP}_\ell^y\rangle\ e^{-i\omega t}\,. \end{aligned} \tag{12}$$

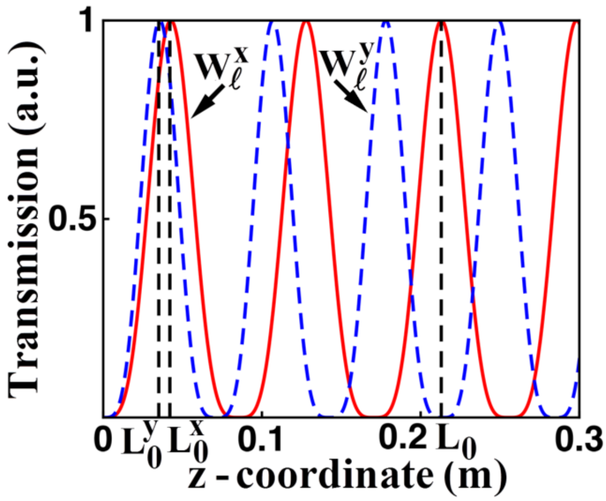

Figure 1. The dependence of the energy W_ℓ^σ stored in the state $|\mathrm{LV}_\ell^\sigma\rangle$ on the fiber's length. The fiber's parameters: $V = 4.16$, $\Delta = 0.001$, $r_0 = 6.3\,\mu m$, the acoustic power $P = 50\,\mathrm{mW}$ and $L_0 = 21$ cm.

This transformation describes a novel type of the optical polarization-dependent mode conversion in fiber acousto-optics based on the particular structure of the modes (2) and non-degenerate propagation constants (4). The key feature of such a conversion is that the topological charge$_\sigma$ ℓ or $-\ell$ (and orbital angular momentum per photon $\hbar\ell$ or $-\hbar\ell$) of the generated field $|\mathrm{LV}_\ell\rangle$ is determined by the direction of linear polarization σ of the incident vortex beam.

The described effect can be useful in developing new acousto-optic devices such, for example, as polarization-controlled optical vortex intensity modulators. Yet, this type of all-fiber wavelength-tunable optical vortex generation and controlling should be especially useful in such vortex-based applications as micromechanics, classical and quantum information encryption, and simulation of quantum computing.

3. References

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Acknowledgments
This work was supported by V.I. Vernadsky Crimean Federal University, grant no. VG 01/2017.