Active quasi-BIC optical vortex generators for ultrafast switching

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Keywords: bound states in the continuum (BIC), optical vortex, ultrafast switching

Abstract

The Pancharatnam–Berry phase induced by the winding topology of polarization around a vortex singularity at bound states in the continuum (BIC) provides a unique approach to optical vortex (OV) generation. The BIC-based OV generators have the potential to outperform their counterparts that rely on spatial variations in terms of design feasibility, fabrication complexity, and robustness. However, given the fact that this class of OV generators originates from the topological property of the photonic bands, their responses are generally fixed and cannot be dynamically altered, which limits their applications to photonic systems. Here, we numerically demonstrate that a silicon photonic crystal slab can be used to realize optically switchable OV generation by simultaneously exploiting the vortex topology in momentum space in conjunction with silicon’s nonlinear dynamics. Picosecond switching of OV beams at near-infrared wavelengths are observed. The demonstrated nontrivial topological nature of the active generators can significantly expand the application of BIC toward ultrafast vortex beam generation, high-capacity optical communication, and mode-division multiplexing.

1. Introduction

Optical beams carrying quantized amounts of orbital angular momentum (OAM), which are featured by the paraxial phase singularities and referred to as optical vortices (OVs) [1], have been extensively studied for their potential in a variety of fields such as optical tweezer [2–4], optical microscopy [5, 6], OAM based communication [7–9], laser [10, 11], quantum communications [12–16], etc. Besides unraveling the intrinsic properties of OVs, the methods of vortex beam generation have attracted tremendous research interest over the past three decades. Approaches such as spiral phase plates [17], computer-generated holograms [18], phase-coded metaatoms [19–21], SAM–OAM coupling based on the Pancharatnam–Berry (PB) phase [16, 22, 23], off-axis generation [24], and so on are proposed to realize the spiral topology. Furthermore, active OV generation with tunable spectral [25–28], temporal [29–31], and OAM [32–35] properties are believed to have the potential to play important roles in numerous applications, including ultrafast vortex beam generation, high-capacity communications and mode-division multiplexing. Nevertheless, most of the reported approaches require a feasible method for implanting a spiral phase-front at the nanoscale as well as the accurate alignment of the devices with a real-space center, which hinders their use in practical applications.

On the other hand, OV generation associated with bound states in the continuum (BIC) [36–42] in two-dimensional (2D) photonic crystal (PhC) slabs offers a flexible alternative approach to realizing high quality OV beams. The corresponding OAM implantation is attributed to the PB phase, which originates from the winding nature of the Bloch modes around the BIC [23]. Relying on the vortex topology in momentum space, this method leads to devices without a real-space center and can also significantly reduce the fabrication complexity. Furthermore, the corresponding polarized modes with winding configuration...
adjacent to the BIC the BIC are featured with a high radiative quality factor (Q-factor) and symmetry-protected robustness, which makes BIC-based OV generators a promising platform for active OAM modulation [43–48]. More importantly, supporting high-Q guided resonance modes, the 2D PhC slabs with a subwavelength thickness simultaneously exhibit strong field confinement, which makes them intrinsically suitable for ultrafast nonlinear switching and less dependent on thermo-optical effects. However, methods for achieving dynamic OV generation from BIC-based photonic systems have rarely been reported.

In addition to the Q-factor, the active response of the material system is the other dominant factor that determines the potential modulation capacity of an optical system when subjected to an external stimulus. High-index dielectric nanostructures showing potential for manipulating light with extremely low loss and high field localization have emerged to be a more practical alternative to their plasmonic counterparts [49–52]. Hydrogenated amorphous silicon (α–Si:H), which is nearly lossless in the near-infrared regime, has proven itself to be an ideal candidate for in situ non-volatile all-optical modulation [53–55] due to its large nonlinearities [56] and fast carrier recombination property [57]. Here, we propose and numerically validate an α–Si:H PhC slab that serves as an effective optically switchable OV generator. The design simultaneously exploits the vortex topology in momentum space and the fast carrier dynamics of α–Si:H. Quantitatively, a correlation factor is introduced to characterize the spatial- and temporal-variation of the modulated OVs. Furthermore, the close relationship between the topological nature of the PhC’ slab and the topological charges of the generated OVs is analyzed based on the Laguerre–Gaussian function and temporal coupling mode theory (TCMT).

2. The PB phase based on winding topologies of resonances

Figure 1(a) illustrates the schematic of the designed OV generator: a 200 nm-thick amorphous silicon (α–Si:H) slab with a square array of cylindrical holes on top of a SiO2 substrate. Unlike the conventional OAM devices [17, 18, 58], with C4-symmetry, this OV generator does not rely on a spatially variant resonance or the winding topology in real space to induce the spiral wavefront. In contrast, its vortex topologies originate from the optical BICs, the bounded states with infinite lifetimes at isolated highly symmetric wave vectors. Importantly, the states of polarization (SOP) of far-field radiation of this type of photonic structure form robust vortex polarization singularities, which reflect the quantized and conserved topological charges [59]. Figure 1(b) shows the band diagram of the structure around a Γ point based on COMSOL eigenmode simulations, in which the color coding represents the logarithmized radiative quality factor (Q-factor) of the Bloch modes. Furthermore, figure 1(c) illustrates the Q-factor of the 2nd-TE-like (referred to as TE2 hereafter) band we focus on here in the vicinity of the Γ point. The lifetime of the resonance goes to infinity at the Γ point, indicating that the system exhibits a symmetry-protected non-radiative BIC. The property of SOP of the TE2 band on an iso-frequency contours (at 1550 nm) is further illustrated in figure 1(d), on a background of the band surface. To clearly depict the origin of the SOP, a comprehensive illustration of the resonance mode at a series of wavevectors in the first quadrant of the momentum space (green double-head arrows) is presented as inserted panels in figure 1(d). In each panel, a vector field of the in-plane electric field is plotted on the color-mapped background of

\[ P = (E_x, E_z) \cdot (k_x, k_z), \]

a value that unambiguously displays the linear polarization perpendicular to the corresponding wavevector (black dash line). Figure 1(d) shows that the SOP winding around the Γ point forms a vortex topology, where the topological charge (+1) can be defined as the winding number of the vortex. This vortex topology reveals that the proposed structure can be viewed as a series of resonators exhibiting \( k_{\text{||}} \)-dependent SOPs [23], which have naturally continuous distribution and are linked to the counterpart OV generators based on spatially variant resonators [19, 60].

When a circularly polarized wave is incident on a spatially variant phase element (for instance, a waveplate with spatially variant local fast axis), the wave at different locations on the element traverses different paths on the Poincare sphere, inducing a spatially variant phase front on its cross-polarized components [61]. The spatially-dependent electric field of the corresponding transmitted wave \( E_o \) can be formulated as follows:

\[
|E_o| = \cos \left( \frac{\Delta}{2} \right) |E_i| - i \sin \left( \frac{\Delta}{2} \right) \left( e^{i\Phi} |R| \langle L \rangle + e^{-i\Phi} |L| \langle R \rangle \right) |E_i|, \tag{1}
\]

where \( |E_i| \) represents the incident electric field, \( \Phi \) represents the spatially variant orientation angle of the fast axis (with respect to the x axis), and \( \Delta \) is the constant phase retardation of the element. To be clear, in this paper the optics convention corresponding to the imaginary unit \( i \) is adopted. Importantly, the element imparts a spatially-dependent phase change of \( 2\sigma\Phi \) on the transmitted waves of the opposite circular polarization, where \( \sigma = \pm 1 \) correspond to the right circularly polarized (RCP) and left circularly polarized (LCP) incident light. Referred to as the geometric (PB) phase, this phase modulation approach gives rise to
Figure 1. Resonance property of the designed α–Si:H/SiO2 OV generator in momentum space near the BIC. (a) Schematic of the device. A 200 nm-thick α–Si:H PhC slab on top of a 130 nm-thick SiO2 substrate ($r = 500$ nm and $a = 1080$ nm). (b) Simulated band diagram around $\Gamma$ point in the $\Gamma$–X and $\Gamma$–M directions. The color represents the logarithmized Q-factor of the resonance modes. (c) Q-factor map of the 2nd-TE-like band. (d) Iso-frequency contour at 1550 nm on top of the eigen frequency map of the 2nd-TE-like band. The double-headed arrows along the contour depict the SOP of the Bloch mode. The insert panels comprehensively illustrate the resonance property of the Bloch mode for a series of wavevectors (dashed black line) in the first quadrant: the black vector field represents the in-plane electric field, and the color mapping denotes the value of $(E_x, E_y) \cdot (k_x, k_y)$ which clearly shows the symmetry property of the mode.

the spiral wavefront observed in BIC-based OV generators, in which SOPs of the Bloch mode winding along the iso-frequency contour exhibit local variation in momentum space (figure 1(d)). In particular, according to the TCMT [23, 62], the transmission process in the proposed PhC slab can be described by the following matrix:

$$T = \begin{pmatrix} t_{ss} & t_{sp} \\ t_{ps} & t_{pp} \end{pmatrix} = t \cdot I - \frac{DD^\dagger}{\gamma + \delta + i(\omega_0 - \omega)} (t + \alpha_z r) \cdot I,$$

(2)

where $t$, $r$, are the scattering coefficient without resonance, $I$ is the identity matrix. $D = (d_s, d_p)^T$ are the coupling coefficients between the resonance and the propagating wave, $\gamma$ and $\delta$ are the radiation loss and the material loss of the resonance respectively, $\omega$ ($\omega_0$) is the angular frequency of the propagating wave (the guided resonance), and $\alpha_z$ is the parity number between the upward and the downward radiation. Clearly, if the SOPs of the guided resonance are nearly perpendicular to the wavevector, i.e., $d_p = 0$, the transmission matrix will add a resonance term, $t_{res}$, to the $s$–pol component as indicated in equation (3).

After a coordinate transformation from an $s$–$p$ basis to a right (left)-handed circularly polarized basis, $|R\rangle - |L\rangle$, the matrix will finally take the form given in equation (4).

$$t_{res} = \frac{|d_s|^2 (t + \alpha_z r)}{\gamma + \delta + i(\omega_0 - \omega)},$$

(3)

$$T = \begin{pmatrix} t_{RR} & t_{RL} \\ t_{LR} & t_{LL} \end{pmatrix} = \left[ t - \frac{1}{2} t_{res} \right] I + \frac{1}{2} t_{res} (e^{i2\phi} |RL\rangle + e^{-i2\phi} |L\rangle \langle R|).$$

(4)

In equation (4), $\phi$ represents the rotation angle (with respect to the $x$ axis) of the wavevector along the iso-frequency contour. The winding nature of the SOP along the iso-frequency contour allows the PhC slab to produce a spiral phase front in $k$-space. From a practical perspective, a slightly divergent circularly polarized (e.g., LCP) Gaussian beam carrying a wide range of $k_l$ components at the resonance frequency can be used to induce the spiral phase front in real space by exploiting the winding topologies of resonances, leading to a RCP OV beam around the vortex topology at the $\Gamma$ point. The corresponding topological charge can be expressed as $l = 2q$, where $q$ is the winding number of the BIC vortex, which, for our design at the TE2 band, is $+1$. High order Laguerre–Gaussian modes can be realized by exploiting BICs in systems of higher symmetry [23]. We note that, primarily relying on the topological properties of BICs,
the OV generation from the proposed OV generator is robust to local imperfections. More importantly, the observed high-Q guided resonances provide the basis for actively controllable vortex generation based on the nonlinear dynamics of α–Si:H as discussed in the later sections. To further evaluate the \( k_r \)-dependent far-field optical response of the proposed structure, full-wave electromagnetic simulations were conducted using COMSOL multiphysics. We note that the scattering simulations are designed based on the eigenmode simulation results (figure 2(a)) along three critical contours, \( \Gamma M, \Gamma X \) and \( k_r = 0.06\pi/a \). Figure 2(a) reveals that because the TE2 band is flat and almost parabolic around the BIC at the \( \Gamma \) point, the iso-frequency contour revealing the high-Q-factor (\( \sim 500 \)) resonance is equivalent to the iso-\( k_r \) contour. The resonance wavelength corresponding to \( k_r = 0.06\pi/a \) is around 1550 nm. Therefore, as the schematic in figure 2(b) shows, the azimuthal angle \( \phi \in [0, \pi/2] \) dependent transmittance spectra of the co- and cross-polarized components of the OV generator under an LCP plane wave illumination were obtained. To study the potential of the proposed device for active vortex generation, two sets of simulations using permittivity of the PhC slab \( \varepsilon_{\text{slab}} = \varepsilon_{\text{Si}} + \Delta \varepsilon \) (state 1, \( \Delta \varepsilon = 0 \), dashed curve) and \( \varepsilon'_{\text{slab}} = \varepsilon_{\text{Si}} + \Delta \varepsilon \) (state 2, \( \Delta \varepsilon = -0.1 + 0.1i \), solid curve), were performed. A similar permittivity modulation realized by exploiting material property dynamics of α–Si under optical excitations has been used to accurately predict the dynamic response from a–Si metasurfaces [53, 54]. Figure 2(c) shows that the transmittance spectra of the cross-pol \( (T_{\text{LR}}) \) components of the OV generator under an LCP plane wave illumination as \( k_r \) moves along the iso-\( k_r \) contour at state 1 (dashed curve) and state 2 (solid curve). At state 1, for each studied incident azimuth angle \( \phi \), a transmittance peak with a narrow bandwidth of \( \sim 15 \) nm is identified at a wavelength around 1550 nm (grey plane). These observed transmittance peaks are closely associated with the high-Q Bloch mode (figure 2(a)). Furthermore, figure 2(d) illustrates the corresponding transmittance spectra of the co-pol \( (T_{\text{LL}}) \) and cross-pol \( (T_{\text{LR}}) \) components when \( \phi = 0 \), in which the dashed (solid) group of curves corresponds to state 1 (state 2). The inset of figure 2(d) provides a close look at the spectra around the resonance wavelength. It
shows, a lowest-order Laguerre–Gaussian (LG) transmitted wave. The metasurface is placed 360\(\lambda\) away from the beam waist. The ultrafast laser pulses at 800 nm are assumed to provide the optical excitation that induces the material property dynamics of \(\alpha\text{–Si:H}\). Simulations were carried out demonstrating OV generation at two states (state 1, \(\Delta\epsilon = 0\), and state 2, \(\Delta\epsilon = -0.1 + 0.1i\)). (b) Electric field intensity, \(|E|^2\), (left) and phase (right) distributions of the generated beam on a plane 640\(\lambda\) away from the slab in state 1. Top row: cross-pol (RCP) component. Bottom row: co-pol (LCP) component. (c) Same as (b), but for state 2.

Figure 3. Dynamic OV generation. (a) Schematic of the active OV generator. A LCP LG\(^{0}_0\) beam \((w_0 = 3\lambda)\) at \(\lambda = 1550\) nm illuminates an \(\alpha\text{–Si:H}/\text{SiO}_2\) PhC slab of finite size \((90 \times 90\) array\), which imparts a spiral phase front on the cross-pol (RCP) transmitted wave. The metasurface is placed 360\(\lambda\) away from the beam waist. The ultrafast laser pulses at 800 nm are assumed to provide the optical excitation that induces the material property dynamics of \(\alpha\text{–Si:H}\). Simulations were carried out demonstrating OV generation at two states (state 1, \(\Delta\epsilon = 0\), and state 2, \(\Delta\epsilon = -0.1 + 0.1i\)). (b) Electric field intensity, \(|E|^2\), (left) and phase (right) distributions of the generated beam on a plane 640\(\lambda\) away from the slab in state 1. Top row: cross-pol (RCP) component. Bottom row: co-pol (LCP) component. (c) Same as (b), but for state 2.

should be noted that the spectra of both \(T_{L2L}\) and \(T_{L2R}\) show a resonance related peak, a phenomenon that is well described by the TCMT summarized in equation (4). Particularly, around the resonance wavelengths, the incident power primarily couples into the in-plane Bloch mode, i.e., the direct transmission \((t\) term in equation (4)) is suppressed, leading to \(T_{L2L}\) decreasing from \(\sim 1\) to \(\sim 0\). Simultaneously, the induced Bloch mode couples out to resonance related transmission, \(t_{res}\), via a radiative channel associated with the radiative loss \((\gamma)\), which, according to equation (4), results in not only the cross-pol \((T_{L2R})\) transmitted signal carrying PB phase, but also the co-pol \((T_{L2L})\) transmitted wave. Therefore, this coupling process gives rise to the observed peaks in both the co-pol and cross-pol transmittance spectra. At state 1, \(T_{L2L}\) and \(T_{L2R}\) with corresponding values around 0.09 and 0.22 are observed, respectively, indicating that the majority of the transmitted power was converted into the cross-pol (RCP) component which carries the PB phase. In sharp contrast, at state 2, \(T_{L2R}\) exhibits a dramatic decrease to \(\sim 0.01\) at 1550 nm while \(T_{L2L}\) slightly increases to 0.12. Figures 2(c) and (d) reveal that in each material state the optical response of the OV generator is insensitive to the azimuthal angle \(\phi\) of the incident wave, which can be attributed to the flattened band surface of the TE\(_2\) band (figure 1(d)) which provides the basis for OV generations via momentum-space polarization vortices centered at BIC. More importantly, the observed largely tunable transmission of the cross-pol components \((T_{L2R})\) essentially offers the potential for actively controllable vortex generation based on optical excitation of the nonlinear dynamics of a–Si:H.

3. Tunable vortex beam generation

To directly study the tunable property of the proposed OV generator, we further simulate the response produced by the proposed PhC slab at the two static states discussed above. As the schematic in figure 3(a) shows, a lowest-order Laguerre–Gaussian (LG\(^0\)) beam at \(\lambda = 1550\) nm (waist radius \(w_0 = 3\lambda\)) was used to illuminate the slab of finite size \((90 \times 90\) array\) that is placed 360\(\lambda\) away from the beam waist. The OV generation switching can be achieved by the pulsed laser excitations which exploit the photoexcited carrier dynamics of \(\alpha\text{–Si:H}\). To investigate the potential of the device for active OV generation, the transmission performance over the aperture was simulated and the results are summarized in figures 3(b) and (c).

It can be seen that, at state 1 \((\Delta\epsilon = \epsilon_{Silicon})\), the LCP Gaussian beam induces the target guided resonance at 1550 nm owing to its plentiful wavevector components on the slab surface, which results in an RCP vortex beam with a spiral phase-front corresponding to the LG\(^0\) mode (figure 3(b)). In particular, a \(4\pi\) spiral phase-front along with the high intensity rings are identified on the phase and field intensity maps (top row in figure 3(b)), respectively, indicating the generation of an OV beam of topological charge \(l = 2\). We note that two concentric high intensity rings are identified in the field intensity \((|E|^2)\) map of the RCP component. In contrast, one single radial peak is predicted from the Laguerre–Gaussian equations [17]. Based on a close inspection of equation (4), we attribute this abnormal feature to the competition between the amount of power coupling into the device and the coupling criteria. As illustrated in the following derivation, this competition phenomenon is the key factor in the design of the OV generator.
As mentioned in the previous section, the excitation of the corresponding in-plane Bloch modes are associated with the horizontal wavevectors offered by the Gaussian beam wavefront. Equation (5) represents the theoretical definition of a Gaussian beam, in which \( w(z) \) is the high-intensity beam radius, \( w_0 = w(0) = \alpha \lambda \) denotes the beam waist, and \( z_0 \) is the Rayleigh range. If the assumption is made that the propagation distance \( z \gg z_0 \), i.e., the Gouy phase term \( \eta(z) \) is negligible, then the second and the third terms in the exponent denote the phase variation of the beam, while the first term describes the exponential decay of the beam intensity along the radial direction. Furthermore, the curvature term, \( k(x^2 + y^2)/(2R(z)) \), which will be the only term that contributes to the horizontal wavevector, contains a radially varying \( k_r \), as seen in equation (6). In particular, at a fixed propagation distance \( z \), \( k_r \) increases along the radial direction.

\[
E = E_0 \frac{w_0}{w(z)} \exp \left( -\frac{x^2 + y^2}{w(z)^2} - ikz - ik \frac{x^2 + y^2}{2R(z)} + i\eta(z) \right),
\]

\[
k_r = k \frac{\sqrt{x^2 + y^2}}{2R(z)} = \frac{kr}{2R(z)}, \quad k_z = k.
\]

Since the BIC at the \( \Gamma \) point ideally prevents the resonance coupling out to far-field radiation, i.e., \( d_s = d_y = 0 \), then the resonance modes used to induce the spiral phasefront must have a non-zero in-plane wavevector, i.e., \( k_{\text{res}} \neq 0 \). To induce these high-Q resonances around the BIC states, the range of \( k_r \) should encompass the value, \( k_{\text{res}} \). As derived in equation (6), \( k_r \) is proportional to the beam radius, \( r \), i.e., the range of \( k_r \) is intrinsically equivalent to the range of \( r \). Hereafter, \( k_{\text{res}} \) along the iso-frequency contour is assumed to be \( 2\pi/a \) (\( a \) is the periodicity), indicating an in-plane wavevector slightly away from the \( \Gamma \)-point. Therefore, we obtain a range of \( k_r \) as defined in equation (7), where a scaling factor \( \beta < 1 \) is used to convert the inequality constraint to an equality constraint. In equation (7), \( k_0 \) is the free space wavevector, and \( \lambda \) is the wavelength of the beam. Since the corresponding range of \( k_r \) is closely related to the in-plane resonance, which primarily determines the couple-in process and the couple-out to OV radiation, it is therefore defined as the coupling criteria.

\[
\frac{k_r}{k_z} = \frac{k_0 \sqrt{x^2 + y^2}}{2R(z)} = \frac{r}{2R(z)} = \frac{r}{2} \frac{z}{z_0} \approx \frac{k_{\text{res}}}{k_0} = \frac{\zeta \lambda}{a} = \beta \frac{r}{2} \frac{z}{z_0}.
\]

On the other hand, the radially exponential decay of the beam intensity limits the input coupling power at a specific beam radius, \( r \). According to equation (7), to reach \( k_{\text{res}} \), the excitation of the resonance is associated with the beam component around \( r = (2\zeta \lambda(z^2 + z_0^2))/\langle az \rangle \). Clearly, the corresponding OV-related incident power, i.e., the input coupling power (at \( r \)) is inevitably lower than that at the beam center. Therefore, the input coupling power at \( r \) is the other limitation of the guided resonance. Quantitatively, the input coupling power is designed to meet the following criteria, \( E_i > e^{-1}E_0 \), which can be further interpreted as equation (8). Substituting the expression developed in equation (7) into equation (8) leads to the inequality given in equation (9) based on the assumption that \( z \gg z_0 \). This inequality, with the factor \( \beta < 1 \), comprehensively illustrates the input coupling power requirement and the coupling criteria during OV modulation, which are regarded as the critical design criteria of the OV modulator. Following these criteria, an incident Gaussian beam (\( \lambda = 1550 \text{ nm} \)) with a waist radius of \( 3\lambda \) is used to excite iso-\( k_r \) Bloch modes (\( k_r = 0.06\pi/a \)) in the \( \alpha-\text{Si:H}/\text{SiO}_2 \) OV generator (\( a = 1080 \text{ nm} \)).

\[
\frac{x^2 + y^2}{w(z)^2} < 1,
\]

\[
\frac{2\zeta}{a_\beta} < \frac{z}{z_0} \frac{1}{\sqrt{2} z^2 + z_0^2} \leq \frac{1}{z_0} \frac{1}{\pi a^2 \lambda} \Rightarrow \frac{2\zeta \pi \alpha \lambda}{a_\beta} < 1.
\]

In short, for the input Gaussian beam, the intensity decreases exponentially along the radial direction from the beam center, while the radial wavevector component, \( k_r \), increases. For the outer annular peaks, the condition \( k_r = k_{\text{res}} \) is met, i.e., the resonance coupling term dominates the direct transmission term in equation (4), \( 1/2t_{\text{res}} \gg t \), which leads to the first maximum on both the cross-pol (RCP) and co-pol (LCP) intensity maps. On the other hand, when \( k_r \) is slightly detuned from \( k_{\text{res}} \) toward zero \( (k_r < k_{\text{res}}, E_i \) approaching \( E_0 \)), resonance coupling no longer dominates the direct transmission. Quantitatively, there is a condition when the resonance-related transmission is equal to the direct transmission, \( 1/2t_{\text{res}} = t \), i.e., when the co-pol component \( (t - 1/2t_{\text{res}}) \) in equation (4) vanishes. This condition exhibits itself as an inner high intensity ring on the cross-pol (RCP) intensity map and a low intensity ring on the co-pol (LCP) intensity map.
On the other hand, at state 2, the RCP wave (top row in figure 3(c)) shows a rather low transmission, although the transmitted beam has an almost identical phase distribution compared with that at state 1. These results indicate the potential of the proposed device for optically switchable OV generation. Furthermore, the corresponding LCP transmitted beam behaves like a $LG_{00}$ beam with no angular phase variation.

It has been learned that compared with an ideal Laguerre–Gaussian vortex beam of the same topological charge, the generated high-order vortex beam behaves as a quasi-Bessel beam that is less divergent [23]. To study the propagation of the generated OV beams in the far-field regime, we performed a Huygens source extension using the simulated electric field on a plane 100 nm away from the exiting surface of the device. From top to bottom panels, figure 4(a) illustrates the beam intensity, phase-front on a plane at $z = 1340 \mu m$ ($z$ denotes the distance to the beam waist) and the side view of the beam ($x = 0$ nm), when $\varepsilon_{\text{slab}} = \varepsilon_{\text{Si}}$ (state 1). A notable annular maximum is observed in the beam profile with a radius of 64 $\mu m$ at $z = 1340 \mu m$.

Figure 4(b) shows the same set of results for state 2. Figures 4(a) and (b) unambiguously reveal the switching effect of OV generation due to the slab’s material property change. As a comparison, figure 4(c) depicts a similar set of results for an ideal $LG_{20}^0$ beam. One should notice that the ring-shaped annular maximum (top panel in figure 4(c)) with the same radius is observed on a plane at $z = 595 \mu m$, indicating much less diffraction broadening of the OV beam generated by the proposed device.

This is made more evident by a comparison between the side view of the beam profiles (bottom row in figure 4). In particular, the generated OV beam shows a divergence angle of 3.4°, while that of the ideal $LG_{00}^0$ beam is 6.3°.
4. Interference and ultrafast switching of VBs

The OAM of the generated OV beam can be further visualized by the interference pattern with a coherent Gaussian beam. Two different interference patterns are derived based on the Huygens extension method. Figure 5(a) illustrate the interference pattern between the OV beam when \( \Delta \varepsilon = 0 \) (state 1) and an ideal RCP Gaussian beam propagating along the x axis. A trident-shaped pattern formed at the singularity is clearly observed, especially in the insert of figure 5(a), which reveals the +2 topological charges at the beam center. As a fingerprint of the high order LG modes, the computer-generated holograms of a similar pattern have been used to produce or down-convert the corresponding Laguerre–Gaussian mode [15, 18, 19, 24, 63]. On the other hand, figure 5(b) shows a dim interference pattern due to the low beam intensity when \( \Delta \varepsilon = -0.1 + 0.1i \) (state 2). Moreover, figure 5(c) illustrates the interference pattern between the generated OV beam and a collinear Gaussian beam at a series of locations along the propagation direction. The spiral phase of two nodes again clearly shows the topological property of the generated OV beam.

Orthogonality is another crucial characteristic of the generated OV beam. According to the theoretical definition, the \( LG_\ell^p \) modes are infinite and ideally orthogonal to each other, which forms a Hilbert space of infinite dimension [12–14]. This feature offers a quantum entanglement freedom of multi-dimension, which is the basis of the secure and the high capacity of quantum communication systems. The normalized orthogonality factor between two ideal LG modes, \(|1\rangle\) and \(|2\rangle\), can be defined as follows,

\[
F_{\text{orth}} = 1 - \frac{\langle 1|2 \rangle}{\langle 1|1 \rangle} = 1 - \frac{\int_0^R \int_0^{2\pi} A_1(r, \Phi) e^{in\Phi} [A_2(r, \Phi) e^{in\Phi}]^* \, dr \, d\Phi}{\int_0^R \int_0^{2\pi} A_1(r, \Phi)^2 \, dr \, d\Phi},
\]

where the amplitude factor, \(A(r, \phi)\), denotes the amplitude and all angularly invariant phase term of the corresponding mode. Clearly, the normalized orthogonality factor approaches to 1 when inner product, \(\langle 1|2 \rangle\), goes to zero, which reflects that the two components’ angularly variant phase terms (\(e^{ni\phi}\) and \(e^{mj\phi}\)) have a non-zero difference, i.e., the topological charges at the beams’ center satisfy the condition \(n - m \neq 0\). In contrast, when \(|2\rangle\) is identical to \(|1\rangle\), the factor will drop to zero. Therefore, \(F_{\text{orth}}\) provides an indicator of the orthogonality between these two modes. Here, we present the spatial and temporal orthogonality dynamics of the OV beams generated from the proposed device where the results are summarized in figure 6.

Figure 6(a) illustrates the \(F_{\text{orth}}\) associated with the two static states of the generator (\(\Delta \varepsilon = 0\) (state 1) and \(\Delta \varepsilon = -0.1 + 0.1i\) (state 2)) as a function of the propagation distance \(z\). The three curves in figure 6(a) represent the orthogonality of three cases, i.e., between two state-1 beams (\(\langle \Delta \varepsilon = 0 | \Delta \varepsilon = 0 \rangle\)), between one state-1 beam and one state-2 beam (\(\langle \Delta \varepsilon = 0 | \Delta \varepsilon = -0.1 + 0.1i \rangle\)), and between two state-2 beams (\(\langle \Delta \varepsilon = -0.1 + 0.1i | \Delta \varepsilon = -0.1 + 0.1i \rangle\)). It can be seen that, for the 1st (red curve) and the 3rd (blue curve) cases, the factor \(F_{\text{orth}}\) is nearly 0 and spatially invariant, illustrating the good self-correlation of the
Figure 6. Spatial and temporal orthogonality dynamics. (a) Spatial orthogonality factor as a function of propagation distance $z$. The results represent the correlation between three pairs of RCP OV beam modes which correspond to the device in the two static states ($\Delta \varepsilon = 0$ and $\Delta \varepsilon = -0.1 + 0.1i$), i.e., $\langle \Delta \varepsilon = 0|\Delta \varepsilon = 0 \rangle$, $\langle \Delta \varepsilon = 0|\Delta \varepsilon = -0.1 + 0.1i \rangle$ and $\langle \Delta \varepsilon = -0.1 + 0.1i|\Delta \varepsilon = -0.1 + 0.1i \rangle$. (b) Temporal orthogonality dynamics between the RCP OV beam modes when the device is optically pumped by an ultrafast laser pulse. A 100 fs laser pulse at 800 nm wavelength with three different pump fluences of $F = 26, 64, \text{and } 135 \mu J cm^{-2}$ are assumed to excite the OV generator. The curves represent the corresponding dynamics on an examination plane located at $z = 1.6 \text{ mm}$, in which the dashed line denotes the long-lasting component primarily determined by the thermo-optical effect in $\alpha$–Si.

Figure 6(b) depicts the temporal orthogonality dynamics. These results are based on the nonlinear model method presented by Della Valle et al. [54] It has been experimentally demonstrated that this theoretical model can accurately describe the nonlinear dynamics of $\alpha$–Si:H nanostructures by comprehensively including all optically induced nonlinear processes (i.e., two-photon absorption, free-carrier injection and relaxation, and lattice heating). In our study, a 100 fs laser pulse at an 800 nm wavelength is assumed to excite the OV generator with three representative pump fluences, $F = 26, 64, \text{and } 135 \mu J cm^{-2}$. These pump fluences range from 2.6% to 13.5% of that used in previous studies [53, 54] and may induce a nonlinear permittivity change $\Delta \varepsilon$ in $\alpha$–Si:H of $-0.02 + 0.02i$ to $-0.1 + 0.1i$. The transient orthogonality factor between the original beam (without pump) and the beam associated with an optical pumping (labeled as $\langle \varepsilon(t) \rangle$) at $z = 1.6 \text{ mm}$ are presented in figure 6(b). The three traces show that, in the first 100 $\text{fs}$ after pump, the orthogonality factor dramatically increases from 0, which corresponds to the self-correlation, to a pump-power-dependent large value, followed by a slower ($\sim 20 \text{ ps}$) recovery process (indicated by the dashed-lines) originating from the lattice heating effect in $\alpha$–Si resonators under pulsed laser excitation. These observations further verify the great potential of the proposed device for ultrafast and optically switchable OV generation.

5. Discussion and conclusions

In summary, we numerically validated an ultrafast nonlinear OV generator with $C_4$-rotational symmetry by simultaneously exploiting the PB phase corresponding to the SOP vortices around BIC and the nonlinear dynamics of the hydrogenated amorphous silicon. A picosecond-scale OAM modulation was theoretically and numerically studied. We note that, since the proposed BIC-based active OV generators are based on the
topological properties of the photonic structure in momentum space, they are intrinsically superior to the real-space devices in terms of feasibility and robustness to local imperfections. Moreover, OVs with different even-order topological charges can be achieved and actively modulated using similar systems which possess different symmetry properties. For instance, OV beams of a topological charge of +6 can be generated from a honeycomb-lattice PhC system with C6 rotational symmetry [23]. This flexibility may further facilitate active up/down conversion between arbitrary LG modes via a proper spectral tuning. By unveiling of the correlation between the topology properties in k-space and the optical spatial modes, the applications of BIC-related photonic devices can be considerably expanded which serves to facilitate the further development of the field of topological photonics.

Acknowledgments

This work was supported in part by the Penn State MRSEC, Center for Nanoscale Science, under Award No. NSF DMR-1420620 and by the Defense Advanced Research Projects Agency (DARPA) (HR00111720032).

Data availability statement

The data that support the findings of this study are available from the corresponding author upon reasonable request.

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