Subwavelength vortical plasmonic lattice solitons

Fangwei Ye,1,3 Dumitru Mihalache,2 Bambi Hu,1,3 and Nicolae C. Panoiu4,*

1Department of Physics, The State Key Laboratory on Fiber Optic Local Area Communication Networks and Advanced Optical Communication Systems, Shanghai Jiao Tong University, Shanghai 200240, China
2Department of Physics, University of Houston, Houston, Texas 77204-5005, USA
3Department of Electronic and Electrical Engineering, University College London, Torrington Place, London WC1E 7JE, UK
4Centre for Nonlinear Studies, and The Beijing-Hong Kong-Singapore Joint Centre for Nonlinear and Complex Systems (Hong Kong), Hong Kong Baptist University, Kowloon Tong, Hong Kong, China

*Corresponding author: n.panoiu@ee.ucl.ac.uk

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We present a theoretical study of vortical plasmonic lattice solitons, which form in two-dimensional arrays of metallic nanowires embedded into nonlinear media with both focusing and defocusing Kerr nonlinearities. Their existence, stability, and subwavelength spatial confinement are investigated in detail. © 2011 Optical Society of America

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When the size of photonic systems is reduced to subwavelength scale, the light confinement and propagation are severely limited by diffraction. Overcoming this major challenge of confining and manipulating optical energy at the nanoscale has attracted a rapidly growing interest in nanophotonics [1]. To this end, one promising approach employs surface plasmon polaritons (SPPs) of metallo-dielectric nanostructures [2]. Although most of the studies of SPPs have focused on their linear optical features, there is a growing interest in the optical properties of SPPs in the nonlinear regime [3–7]. This is primarily because the large increase of the optical field induced by SPPs can significantly enhance nonlinear optical effects, as well as the promising applications of SPPs to active nanodevices. A new approach to achieve subwavelength confinement and control of the optical field, employing either one-dimensional (1D) or two-dimensional (2D) arrays of metallic nanowires embedded in an optical medium with Kerr nonlinearity, was recently proposed [6]. Under certain conditions the optical nonlinearity induced by the field of SPPs exactly balances the discrete diffraction in the plasmonic array, and as a result both 1D and 2D fundamental (vorticityless) plasmonic lattice solitons (PLSs) are formed in arrays of metallic nanowires [6]. Because the radius, a, of the plasmonic nanowires and their separation distance, d, are much smaller than the operating wavelength, λ, the spatial width of such PLSs can be significantly smaller than λ.

In this Letter, we study the unique features of nonlinear plasmonic vortices in 2D arrays of metallic nanowires embedded into a Kerr medium (Fig. 1). In particular, we demonstrate, to the best of our knowledge for the first time, that vortical PLSs of subwavelength extent in the transverse section exist in this 2D geometry. Thus, in addition to the fundamental (vorticityless) PLSs, which form in both 1D and 2D arrays, 2D plasmonic arrays support vortical solitons, i.e., solitons with a topological phase change of 2π along a closed contour around soliton’s phase singularity. Similar to the case of fundamental PLSs, (subwavelength) vortical PLSs are found to exist in both focusing and defocusing Kerr media.

Our analysis is based on the coupled mode theory for discrete solitons [8], which was developed in [6] for coupled metallic nanowires. To briefly outline the theoretical approach, the total electric and magnetic fields are expanded as a superposition of the modes of a single nanowire (assumed to have only the fundamental TM mode), and using the conjugated form of the Lorentz reciprocity theorem one obtains

\[
i \frac{d\phi_{m,n}}{dz} + \kappa (\phi_{m+1,n} + \phi_{m-1,n} + \phi_{m,n+1} + \phi_{m,n-1}) + \mu (\phi_{m+1,n+1} + \phi_{m-1,n-1} + \phi_{m,n+1} + \phi_{m,n-1}) + \gamma |\phi_{m,n}|^2 \phi_{m,n} = 0.
\]

Fig. 1. (Color online) Top panels, schematic of a 2D square array of nanowires and the band predicted by the coupled mode theory. Bottom panel, the normalized coupling constant versus separation distance, determined by using the coupled mode theory and exact numerical simulations. The radius of the nanowires is a = 40 nm.

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where \( z \) is the longitudinal coordinate, \( \phi_{m,n} \) is the normalized mode amplitude in nanowire \((m, n)\), \( \kappa \) and \( \mu \) are the coupling coefficients between neighboring and next-neighboring nanowires, respectively, and \( y \) is the normalized nonlinearity coefficient. Note that \( y > 0 \) \((y < 0)\) for a medium with self-focusing (self-defocusing) Kerr response. We emphasize that, unlike in the case of conventional dielectric waveguides, for our plasmonic array, \( \kappa < 0 \), \( \mu < 0 \), so that the linear dispersion, \( k_z = 2 \kappa [\cos(k_x d) + \cos(k_y d)] + 4 \mu \cos(k_x d) \cos(k_y d) \), forms a concave surface with anomalous diffraction occurring at the \( \Gamma \) point \((k_x = k_y = 0)\), while normal diffraction occurs at the \( M \) points \((k_x = k_y = \pm \pi/d)\) (Fig. 1). We note that a negative coupling coefficient also arises in all-dielectric Bragg waveguides [9–12] where the coupling mechanism is a radiative one; this is in contrast to the plasmonic arrays [4–6, 13] where the mode coupling is a purely evanescent one. Typical values of the coupling coefficients are \( \kappa = -9.02 \times 10^4 \text{m}^{-1} \) and \( \mu = -1.48 \times 10^4 \text{m}^{-1} \) and correspond to \( a = 40 \text{nm} \) and \( d = 8a \). Note that \(|\kappa|\) is almost an order of magnitude larger than \(|\mu|\), which illustrates the steep decrease of the field away from the nanowires. We have also determined \( \kappa \) by using a software tool (COMSOL), which solves the full set of Maxwell equations, the conclusion being that the coupled mode theory provides an excellent description of the mode coupling for separation distance as small as \( d \approx 3.5a \) (see Fig. 1).

The soliton solutions to Eq. (1) are sought in the form \( \phi_{m,n}(z) = u_{m,n} \exp(i\beta z) \), where the amplitudes \( u_{m,n} \) are independent of \( z \) and \( \beta \) is the soliton wave number (normalized by \( \kappa \)). The soliton is characterized by its power \( P = \sum_{m,n} |\phi_{m,n}|^2 \). This ansatz is inserted into Eq. (1) and the resulting system is solved numerically. The amplitudes are then used to reconstruct the fields of the PLSs. To quantify the transverse size of solitons we use the effective radius, \( R = \left( \iint |E|^2 (x - x_0)^2 + (y - y_0)^2 \right)^{1/2} dx dy / \iint |E|^2 dx dy \) (see Fig. 3), where \( (x_0, y_0) = \iint (x, y) |E|^2 dx dy \) is the mean center position. Also, we assumed that the nanowires are made of Ag and used the Drude model, \( \epsilon_m(\omega) = 1 - \frac{\omega_p^2}{\omega(\omega + i\gamma)} \), for the permittivity of the metal. For Ag, \( \omega_p = 13.7 \times 10^{15} \text{rad/s} \) and \( \gamma = 2.7 \times 10^{13} \text{rad/s} \). For simplicity, we assume in the following that metal is lossless. However, when the optical losses are included, the absorption coefficient is about \( 900 \text{cm}^{-1} \), which corresponds to a decay length of about \( 11 \mu m \). Thus, a truly solitonlike propagation would require a gain of about \( g = 900 \text{cm}^{-1} \), which can be easily achieved experimentally [14].

We recall that both 1D and 2D fundamental PLSs can be of two types, i.e., either staggered or unstaggered [6]. In the case of unstaggered (staggered) PLSs, the phase difference of the mode amplitude in adjacent nanowires is equal to zero (\( \pi \)). Importantly, they exist and are stable only if their power exceeds some threshold. Also, due to the inverted linear dispersion relation, staggered (unstaggered) solitons are formed in self-focusing (self-defocusing) media, which is opposite to the case of dielectric waveguide arrays (see [8]).

In addition to fundamental PLSs, 2D plasmonic arrays also support a new type of PLSs, which has no counterpart in 1D geometries, i.e., vortical PLSs; these belong to the class of nonlinear discrete vortices [15].

\[ \text{(a) } \beta = 3.82, \delta m_{nl} = 0.11, \text{ (c) } \beta = 5.2, \delta m_{nl} = -0.18. \]

\[ \text{Radius } R \text{ (e) and power } P \text{ (f) versus } \beta. \] The solid and dashed curves stand for \( d = 7a \) and \( d = 8a \), respectively. In the inset, \( \delta m_{nl} \) versus \( \beta \) for \( d = 8a \), in focusing medium. In all plots, \( \lambda = 1550 \text{nm} \) and \( a = 40 \text{ nm} \).

\[ \text{(a) } \beta = 3.82 \text{ (top panels), showing the intensity profile at } z = 0 \text{ and at } z = 90 \mu m. \]

\[ \text{Evolution of maximum electric field amplitude with propagation distance for the unstable vortical PLS (bottom panel).} \] In all plots, \( \lambda = 1550 \text{nm} \) and \( a = 40 \text{ nm} \).
we focus on off-site vortical PLSs, i.e., compact vortex states whose singularity is located between lattice sites. Similarly to fundamental PLSs, vortical PLSs exist both in self-focusing [Figs. 2(a) and 2(b)] and self-defocusing media [Figs. 2(c) and 2(d)]. As can be easily seen from the phase profiles, the vortex has a topological charge equal to 1 (2D PLSs with topological charge 2 also exist, although not shown here). The phase profile for PLSs in focusing media features a “staggered” pattern at the location of the plasmonic nanowires, which is a typical feature of “gap” vortices. However, one should mention that the staggered vortical PLS resides at the semi-infinite gap (above the band) and, thus, is not a “gap” vortex. This is again a consequence of the inverted linear dispersion band. The properties of vortical PLSs are summarized in Figs. 2(e) and 2(f). Several differences from fundamental PLSs are to be mentioned. First, unlike the fundamental PLSs where \( \beta \) can be very close to the band edges, vortical PLSs can only exist in the gaps at a finite distance from the band edges. This is due to the fact that the vortex mode differs significantly from a Bloch mode of the linear plasmonic array. Second, as vortical PLSs possess four intensity maxima, they always have a larger width as compared to that of fundamental PLSs (under the same change of nonlinear refractive index). Therefore, forming a subwavelength vortex generally requires a stronger nonlinearity. Nevertheless, we find that subwavelength vortical PLSs can be achieved under experimentally accessible conditions. For example, the vortex presented in Figs. 2(a) and 2(b) has a radius \( R = 0.3 \lambda \) requiring a nonlinear change of refractive index \( \delta n_{\text{val}} = 0.11 \). Note that an index change of \( \sim 0.14 \) was reported in [17]. Further, the size of vortical PLSs can be significantly reduced if the wavelength is scaled down, thus, the requirement for a strong nonlinear change of the refractive index can be relaxed. Finally, we mention that, as expected, vortical PLSs require a threshold power for their formation.

A relevant issue associated with the families of vortical PLSs is their stability. To address this issue, we integrate Eq. (1) by using the Runge–Kutta method with the input condition being the stationary soliton solution, to which random noise is added. Our extensive numerical simulations show that the vortical PLSs can be stable in certain domains of their existence regions, although their stability domains are more limited in comparison with those of vorticityless PLSs. The stability region of the vortical solitons is characterized by the dependence \( P = P(\beta) \) [see Fig. 2(f)]. Specifically, unlike vorticityless PLSs, which change from unstable into stable solitons exactly at the point where \( dP/d\beta \) changes sign, vortical PLSs remain unstable until \( |\beta| \) exceeds some value. A typical unstable propagation of a vortical PLS is shown in Fig. 3. The vortex soliton quickly loses its screwed phase structure and decays into a (linear) mode of one of the nanowires. Note that the stability domain of vortical PLSs increases if we operate at a blue-shifted wavelength.

In conclusion, we presented the first study of families of vortical plasmonic lattice solitons in 2D arrays of metallic nanowires embedded into a host medium with Kerr nonlinearity. Families of stable subwavelength plasmonic vortex solitons are found to exist in both focusing and defocusing Kerr media. We expect that due to the subwavelength character of the vortical PLSs, our results will foster exciting new applications in subwavelength nanophotonics and plasmonics, including ultrasmall detectors and active nanodevices.

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References