

# Observation of surface gap solitons in semi-infinite waveguide arrays

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We report on the observation of *surface gap solitons* found to exist at the interface between uniform and periodic dielectric media with defocusing nonlinearity. We demonstrate strong self-trapping at the edge of a LiNbO<sub>3</sub> waveguide array and the formation of staggered surface solitons with propagation constant inside the first photonic band gap. We study the crossover between linear repulsion and nonlinear attraction at the surface, revealing the mechanism of nonlinearity-mediated stabilization of the surface gap modes.

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Interfaces between different physical media can support a special class of localized waves known as surface waves. Surface waves attract great attention with their possible application in surface sensing and probing. In periodic systems, *staggered surface waves* are often referred to as Tamm states [1], first identified as localized electronic states at the edge of a truncated periodic potential. Direct observation of electronic surface waves in natural crystalline materials remains beyond practical reach, but successful efforts were made to demonstrate their existence in nanoengineered superlattices [2]. In optics, linear staggered surface modes or Tamm states have been demonstrated at surfaces separating periodic and homogeneous dielectric media [3].

The nonlinear response of materials makes possible the dynamic control of surface localization. This ability has generated great interest in the study of nonlinear surface waves in different fields of physics and most extensively in optics. A self-focusing optical nonlinearity enables the existence of localized waves at interfaces between homogeneous dielectric media [4–6] where no linear modes exist. However, such surface modes are typically associated with high power requirements, and are not possible at all if the nonlinear response is self-defocusing. The combination of periodicity and nonlinearity allows to overcome both of these limitations due to the ability of periodic structures to dramatically modify beam diffraction. This leads to a wealth of different types of modes localized at and near the surface [7]. Furthermore, nonlinear surface states hold promise for optical switching applications utilizing the mechanism of low power beam delocalization and higher power confinement at the surface.

Self-trapping of light near the boundary of a *self-focusing* photonic lattice was predicted recently [8] and demonstrated in experiment [9] through the formation of *discrete surface solitons* at the edge of a nonlinear

waveguide array. In *self-defocusing* materials the existence of *surface gap solitons* at the interface between uniform and periodic media was also recently predicted theoretically [8, 10]. In this case, light localization occurs inside the photonic bandgap in the form of staggered surface modes [10]. This enables one to extend the analogy with the localized electronic Tamm states into the nonlinear regime, so the surface gap solitons can be termed as *nonlinear Tamm states*.

In this Letter we study experimentally self-action of a narrow beam propagating near the edge of a LiNbO<sub>3</sub> waveguide array with defocusing nonlinearity. For the first time to our knowledge, we observe the formation of surface gap solitons, or nonlinear Tamm states. While linear surface modes *do not exist* in this type of system, discrete light self-trapping is observed in the nonlinear regime above a certain threshold power when the propagation constant is shifted into the gap of the photonic transmission spectrum. By employing a simple nonlinear discrete model [11], we describe the crossover from discrete diffraction and surface repulsion in the linear regime, to the appearance of a purely nonlinear localized surface state at higher optical intensities. We discuss the physical mechanism of the nonlinearity-induced stabilization of the staggered surface modes.

In our experiments, we study nonlinear surface localization in a semi-infinite array of single-mode optical waveguides fabricated by a Titanium in-diffusion process in a mono-crystal lithium niobate (LiNbO<sub>3</sub>) wafer, similar to that recently used for the observation of discrete gap solitons [12, 13]. The fabrication process, described in Ref. [13], results in a high-quality waveguide array with refractive index contrast  $\Delta n = 3 \times 10^{-4}$ , waveguide spacing  $d = 9.0 \mu\text{m}$ , sample length 50 mm, and a total of 100 waveguides. Inset in Fig. 1(b) shows schematically the geometry of the waveguide array. The LiNbO<sub>3</sub> sam-

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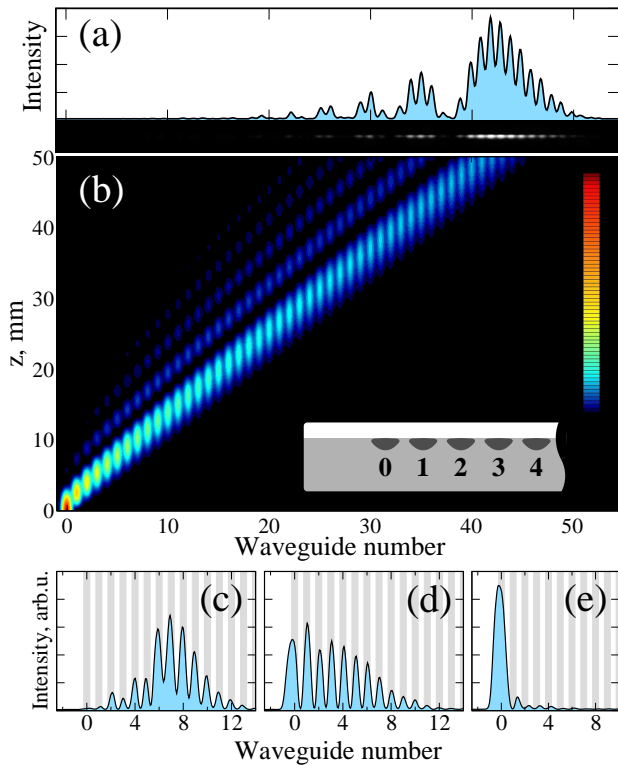


FIG. 1: (color online) Linear propagation of a narrow low-power beam when only the edge waveguide of the array is excited. (a) Measured transverse output intensity profile ( $P = 0.1 \mu\text{W}$ ) and (b) corresponding theoretically calculated longitudinal propagation inside the sample. Inset in (b) shows the waveguide geometry. (c-e) Formation of the surface gap soliton at the array output 920, 1050, and 1550 s, respectively, after the input beam power is increased to  $P = 0.5 \text{ mW}$ . Grey shading marks the waveguide positions.

ple exhibits a strong photovoltaic effect which leads to defocusing nonlinearity at visible wavelengths.

In the experimental setup an extraordinarily polarized probe beam from a cw Nd:YVO<sub>4</sub> laser ( $\lambda = 532 \text{ nm}$ ) is focused by a microscope objective ( $\times 20$ ) to a full width at half-maximum (FWHM) of  $2.7 \mu\text{m}$  at the input face of the sample, and injected into the waveguide at the edge of the array. The propagated wavepacket at the output of the sample is imaged onto a CCD camera. The FWHM of the individual waveguide mode is  $6 \mu\text{m}$  and  $3 \mu\text{m}$  in horizontal and vertical directions, respectively, allowing for a single-waveguide input coupling. The waveguide array is externally illuminated by a white-light source in order to control the nonlinear response time. As shown in Ref. [13] single site excitation provides an efficient method for excitation of gap solitons in periodic defocusing nonlinear materials, provided the refractive index contrast exceeds a certain threshold. In this case the periodic structure appears equivalent to a discrete system [13] and can be well described by a nonlinear discrete model.

At low laser power ( $0.1 \mu\text{W}$ ), we observe two ma-

ior effects. First, due to coupling between neighboring waveguides the probe beam experiences discrete diffraction and spreads out in the horizontal plane upon propagation. Second, the beam shifts dramatically to the right indicating a strong repulsive effect of the surface. Figure 1(a) shows the experimental image of the output intensity distribution and the corresponding transverse intensity profile. After linear propagation through the array the beam profile acquires a complex form, where the major lobe is centered approximately 42 lattice sites away from the input excitation point ( $n = 0$  at the edge of the array) due to the surface repulsion. Figure 1(b) shows the corresponding optical intensity distribution inside the sample, calculated with the help of a simple analytical formula derived from a discrete model  $a_n(z\kappa) = A_0 i^n [J_n(2z\kappa) + J_{n+2}(2z\kappa)]$ , where  $a_n(z\kappa)$  is the discrete mode amplitude in the  $n$ -th waveguide,  $A_0$  is the initial field amplitude in the input waveguide  $n = 0$ ,  $z$  is the propagation distance, and  $\kappa$  is the intersite coupling coefficient [8]. In Fig. 1(b) the discrete mode amplitudes have been multiplied by the continuous waveguide mode intensity profile, and the agreement with the experimental observation is found to be excellent. The coupling coefficient is estimated to be  $\kappa = 0.46 \text{ mm}^{-1}$ , implying a total propagation of 23 coupling lengths.

Increasing the laser power leads to spatial beam self-action through the defocusing photovoltaic nonlinearity. The slow response of the nonlinearity allows us to monitor directly the transient temporal dynamics of self-trapping and soliton formation. Figures 1(c-e) show the output beam intensity profile at times 920, 1050, and 1550 s, respectively, after the beam power is increased to  $P = 0.5 \text{ mW}$ . The wavepacket is seen first to contract and shift towards the edge of the array, indicating a nonlinearity-induced suppression of the surface repulsion [Fig. 1(c)]. Then partial self-trapping at the surface occurs, with a tail of intensity lobes extending into the periodic structure [Fig. 1(d)]. A series of zero intensity points between these lobes indicates the self-induced dynamic formation of a *staggered phase structure* which is clearly absent in Fig. 1(a). Eventually, a strongly localized *surface gap soliton* is formed [Fig. 1(e)]. The asymmetry of the photonic structure is reflected in the shape of the trapped beam which decays monotonically into the continuum while showing damped oscillations inside the array, resembling the structure of a *truncated Bloch mode*. The defocusing nonlinearity effectively decreases the contrast of the surface waveguide, causing the localized mode to broaden and penetrate substantially into the continuous medium.

In order to study in detail the crossover between linear diffraction and nonlinear self-localization, we measure the surface gap soliton formation time as a function of the probe beam power. As seen in Fig. 2 the formation time increases dramatically for decreasing input power until, below a certain critical power, no lo-

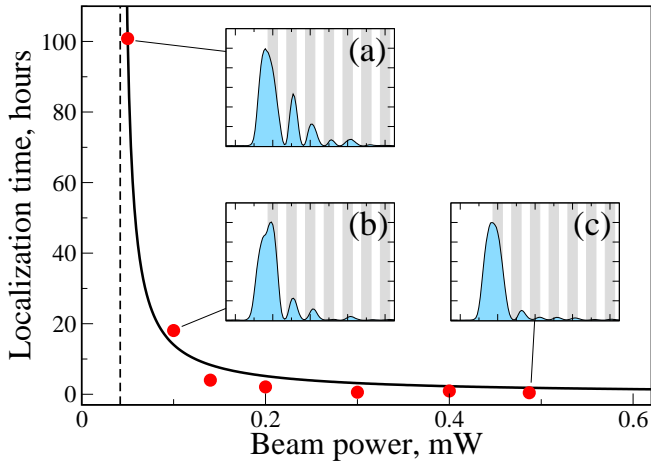


FIG. 2: (color online) Measured surface localization time vs. probe beam power. Solid curve:  $A + B/(P - P_{\text{th}})$  fit to experimental data (red dots). Vertical dashed line marks the threshold power ( $P_{\text{th}} = 0.042$  mW). (a-c) Beam intensity profiles of decreasing width corresponding to the indicated points.

calized surface mode is observed. This observed critical slowing down indicates the existence of a threshold power below which the nonlinear response is too weak to cause self-trapping. While the analytical form of the dynamics of the soliton formation near the threshold is not known, the value of the threshold power was estimated as  $P_{\text{th}} = 42 \mu\text{W}$  by modeling the measured dependence of the soliton formation time (Fig. 2, red dots) on the beam power with a simple singular relation  $A + B/(P - P_{\text{th}})$  (Fig. 2, solid curve). Figures 2(a-c) show the beam intensity profiles corresponding to the indicated data points. The width of the localized mode decreases for increasing beam power, spanning about three lattice sites immediately above threshold [Fig. 2(a)], and approximately a single lattice site at higher power [Fig. 2(c)]. The decrease of the beam width is due to the fact that stronger beam self-action at higher power leads to a deeper surface defect, and hence more pronounced beam localization.

An essential and unique feature of the observed surface gap solitons is the *staggered phase structure* of the beam tail inside the periodic medium. The alternating phase of the field lobes reflects the fact that the propagation constant of the self-localized mode lies within the photonic bandgap at the edge of the first Brillouin zone. To verify that this is indeed the case in the experiment, we interfere the output beam with a vertically inclined plane reference wave. Figure 3(a) depicts a three-dimensional representation of the spatial beam intensity distribution of the broad surface gap soliton observed near the threshold [Fig. 2(a)]. Figure 3(b) shows a two-dimensional intensity plot of the associated interference pattern. A half-period vertical shift of the interference fringes, corresponding to an exact  $\pi$  phase jump in the horizontal beam direction, is clearly observed between each pair of lobes in the struc-

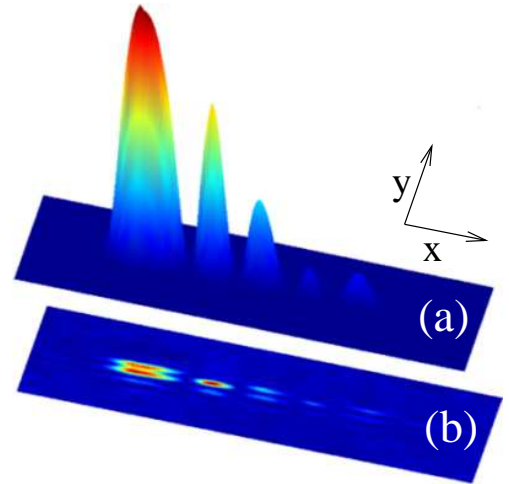


FIG. 3: (color online) (a) Three-dimensional representation of the surface gap soliton observed near the threshold [Fig. 2(a)]. (b) Plane-wave interferogram demonstrating the staggered phase structure of the surface gap soliton.

ture [Fig. 3(b)]. The phase is seen to be constant in the continuous region. The staggered phase structure inside the array and the plane phase in the continuum are signatures of the two different localization mechanisms in play [10]. The mode is confined from the continuum by total internal reflection, while Bragg reflection is responsible for localization inside the periodic structure.

In order to get a deeper insight into the physics of staggered surface soliton formation, we consider the system of coupled-mode equations [14] for the normalized mode amplitudes  $E_0$  and  $E_n$  ( $n = 1, 2, \dots$ ), assuming weak coupling between the neighboring waveguides,

$$\begin{aligned} i \frac{dE_0}{dz} + E_1 + \mathcal{F}(E_0)E_0 &= 0, \\ i \frac{dE_n}{dz} + (E_{n+1} + E_{n-1}) + \mathcal{F}(E_n)E_n &= 0, \end{aligned} \quad (1)$$

where  $\mathcal{F}(E) = \gamma/(1 + |E|^2)$  ( $\gamma > 0$  for defocusing nonlinearity) accounts for the saturable character of the photovoltaic nonlinearity [15] and allows for correct description of the observed effects. This treatment is complimentary to Ref. [10] where pure Kerr nonlinearity was considered.

Looking for stationary solutions in the form  $E_n(z) = \exp(i\beta z)E_n$ , we obtain the linear spectrum of extended modes,  $\beta = 2 \cos k$ , ( $0 \leq k \leq \pi$ ). No localized surface mode exists in the linear regime, as this would require large index contrast between the waveguides and the continuum. However, the presence of defocusing nonlinearity in the model (1) can give rise to localized states. To find them we solve numerically the corresponding stationary equations by a multi-dimensional Newton-Raphson scheme. Since we are interested in surface localized modes, we look for states with maxima near the surface which decay quickly away from the edge of the array. This approach is similar to the earlier studied cases

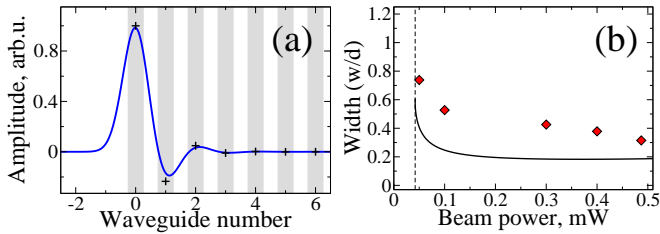


FIG. 4: (a) Example of the staggered surface mode calculated by use of the discrete nonlinear model (1) at  $\gamma = 8$ . The discrete mode amplitudes are marked by + signs. (b) Normalized width of the localized surface state calculated numerically (solid curve) and measured experimentally (diamonds) as a function of the beam power. Vertical dashed line marks the threshold power ( $P_{th} = 0.042$  mW).

of Kerr nonlinearity [8, 10] and agrees qualitatively with the analysis of the continuous model [10].

Figure 4(a) shows the surface mode calculated by use of the discrete nonlinear model (1) and multiplied by the waveguide mode-field profiles. Using the discrete model (1) provides a reasonable agreement with the experimental data. An example of such a comparison is shown in Fig. 4(b) for the width of the localized surface state calculated numerically (solid curve) and measured experimentally (diamonds) as a function of the beam power. Nonlinear coupling constant  $\gamma$  and scaling of the propagation variable  $z$  were adjusted such that the model yields the same diffractive field distribution (Fig. 1) and critical power ( $P_{th}$ ) as the experiment.

Despite being approximate, the discrete model can be employed to reveal the physical mechanism of the nonlinearity-induced surface mode stabilization. To this end we follow earlier studies [16, 17] and calculate the effective energy of the mode,  $H = -\sum_n \{(E_n E_{n+1}^* + E_n^* E_{n+1}) + \gamma \ln(1 + |E_n|^2)\}$ , as a function of its collective coordinate  $X = P^{-1} \sum_n n |E_n|^2$ , where  $P = \sum_n |E_n|^2$  is the mode power. We employ a numerical constraint method [17] which enables to calculate all possible solutions of Eqs. (1) for fixed values of power  $P$ . In this way one can find solutions centered not only at particular waveguides but also intermediate states centered between waveguides. Applying this technique we constructed the effective potential of the surface mode in a semi-infinite array. Figure 5 show this effective potential,  $U_{eff}(X) \equiv -H(X)$ , calculated for two different power values. The extremal points of this curve defined by the condition  $dH/dX = 0$  correspond to stationary localized solutions. In comparison with an infinite array, the truncation of the waveguide array introduces an effective *repulsive surface potential*, which is combined with the periodic potential of the array. As a result, discrete surface modes are not possible in the linear regime. As we see from Fig. 5(a), for low powers there exists no solution of the equation  $dH/dX = 0$  at the surface site  $n = 0$ , and the surface repels the input beam as clearly

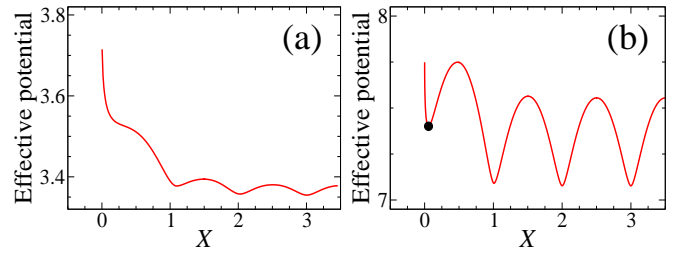


FIG. 5: (color online) (a,b) Effective potential of the localized gap modes vs. the collective coordinate  $X$  below and above the threshold power, respectively. Integer values of  $X$  correspond to the waveguide numbers. Black dot in (b) refers to the stationary solution shown in Fig. 4(a).

observed in experiment [see Fig. 1(a,b)]. However, when the input power exceeds the threshold value, discreteness overcomes the surface repulsive force and the localized state at  $n = 0$  in the form of a surface gap soliton becomes possible [Fig. 5(b)].

In conclusion, we have shown theoretically and demonstrated experimentally that gap solitons can be stabilized near the surface of a periodic medium with self-defocusing nonlinearity in the form of staggered surface modes, providing the first experimental evidence of a nonlinear analog of surface Tamm states in optics.

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