Large phase shift of spatial solitons in lead glass

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The phenomenon of the large phase shift of the strongly nonlocal spatial optical soliton was predicted by Guo et al., within the phenomenological framework [Q. Guo, el al., Phys. Rev. E 69, 016602 (2004), but has not been experimentally confirmed so far. We theoretically and experimentally investigate the large phase shift of that propagating in the lead glass. It is verified that the change of the optical power carried by the optical beam about 10 mW around the critical power for the soliton can lead to a \(\pi\) phase shift, which would be of its potential in the application of all-optical switchings. © 2013 Optical Society of America

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The extensively investigated strongly nonlocal spatial optical solitons (SNSOSs) are first predicted by Snyder and Mitchell [1]. Compared with their local counterparts, SNSOSs can take on complex forms, such as high order solitons [2–5] and even incoherent solitons [6–8]. More importantly, the phase shift of the SNSOS is quite large [9]. This is an essential attribute of the SNSOSs but ignored by Snyder and Mitchell [1].

Assuming that the scalar field of the monochromatic light is

\[ E(x, y, z, t) = A(x, y, z) \exp[-i(\omega t - kz)], \]

\(A\) is the paraxial optical beam, \(k = \omega n_0/c\), and \(n_0\) is the linear refractive index. \(kz\) is a linear phase shift after the propagating distance \(z\) which can be called the geometrical phase shift, while the argument of the paraxial optical beam, \(\arg A\), is “the additional phase shift” that will be abbreviated to “the phase shift” in the following. In the linear case, the increase rate of the phase shift per unit distance is far slower than that of the geometrical phase shift [10]. This is the reason why the phase shift is not treasured all along. Even in the nonlinear case, the phase shift per unit distance of the local soliton was found to be \(1/(2kw_0^2)\) \[11\], where \(w_0\) is the soliton width, which is the same order with the result for the linear optical beam.

Based on the phenomenological Gaussian response function, Guo et al., predicted that the phase shift rate per unit propagation distance is \((w_m/w_0)^2/(2kw_0^2)\) for the SNSOSs [9], where \(w_m\) is the width of the response function. The phase shift is much larger than the local case since the strong nonlocality means \(w_m > 10w_0\) at least. The phase shift rate in the nematic liquid crystal, the first-found material with the strongly nonlocal nonlinearity [12], was found to be \(\pi^{1/2}(w_m/w_0)/(2kw_0^2)\) \[13\]. Though this result is slower than that obtained based on the phenomenological model, it is 10 times faster at least than the results for the local soliton and the linear beam. It was also pointed out [14] that \(\pi\) phase shift of the signal SNSOS can be obtained within a very short given distance via adjusting the pump SNSOS power with the aid of the cross modulation between the SNSOSs. Because, however, the additional phase shift of the SNSOS is completely covered up by the geometrical phase shift during their propagation, it is somewhat difficult to experimentally demonstrate the large phase shift of the SNSOS. Someone even doubted that whether the conclusion of the large phase shift would be right [15].

In this Letter we theoretically and experimentally investigate the large phase shift of the SNSOSs in lead glass. Based on the principle of Mach-Zehnder interferometer, we test and verify the linear modulation of the SNSOS phase by the power of itself. A \(\pi\) phase shift is obtained by changing the soliton power about 10 mw around the critical power, which demonstrates a high modulation sensitivity.

The medium we concern is the lead glass with an extremely large range of nonlocality of the thermal self-focusing type nonlinearity [2, 3, 8, 16]. The propagation behavior of the light beam in this system is governed by the coupled equations [2, 16], which are expressed in the cylindrical coordinate system \((R, \phi, Z)\) for \(Z\)-axis symmetric geometry

\[ 2ik \frac{\partial A}{\partial Z} + \frac{1}{R} \frac{\partial}{\partial R}(R \frac{\partial A}{\partial R}) + 2k^2 \frac{\Delta n}{n_0} A = 0, \quad (1a) \]

\[ 1 \frac{\partial}{\partial R}(R \frac{\partial T}{\partial R}) = -\frac{\alpha}{\kappa} I(R), \quad (1b) \]

where \(\alpha\) and \(\kappa\) are respectively the absorption coefficient and the thermal conductivity, \(\Delta n = \beta \Delta T\) with \(\beta\) being the thermo-optical coefficient, and \(I(R) = |A(R)|^2\). We rewrite Eq. (1) in a dimensionless form:

\[ i \frac{\partial a}{\partial z} + \frac{1}{r} \frac{\partial}{\partial r}(r \frac{\partial a}{\partial r}) + Na = 0, \quad (2a) \]

\[ 1 \frac{\partial}{\partial r}(r \frac{\partial N}{\partial r}) = -|a|^2, \quad (2b) \]

where \(r = R/w_0, z = Z/(2kw_0^2), a = A/A_0, A_0^2 = n_0\kappa/(2\alpha\beta k^2 w_0^4)\) and \(N = 2k^2 w_0^2 \Delta n/n_0\).

Following the Snyder’s method [1], the nonlinear index is expanded and only kept the first two terms of Taylor series:

\[ N = N^{(0)} - r^2 N^{(2)}. \quad (3) \]
Assuming the beam of the Gaussian function form

\[ a = \frac{\sqrt{P_0}}{\sqrt{2\pi} w(z)} \exp[i\theta(z)]\exp(-\frac{a^2}{2w(z)^2}) \tag{4} \]

where \( P_0 = \int |a(x', x, y')|^2 dx' dy' \) is the normalized light power, we can obtain \( N^{(0)} \) by directly integrating Eq. (2b) twice

\[ N^{(0)} = \frac{P_0}{4\pi} \left[ \Gamma(0, \frac{R_0^2}{w^2(z)}) + \ln\left(\frac{R_0^2}{w^2(z)} + \gamma\right) \right], \tag{5} \]

where \( \gamma \) is Euler’s constant which equals to 0.5772156649, and \( R_0 \) is the diameter of the cross section of the lead glass.

\( \theta \) in Eq. (4) is just the phase shift of the beam. We rewrite it into two terms according to the two terms of the nonlinear index (in Eq. (3)):

\[ \theta = \theta^{(0)} + \theta^{(2)}, \tag{6} \]

where \( \theta^{(0)} = N^{(0)} z \) is the zero-order term of the phase shift. By the method in Ref. [9], inserting Eq. (3) and Eq. (4) into Eq. (2b), the beam width and the second-order term of the phase shift can be obtained

\[ w(z) = \sigma + (1 - \sigma) \cos(bz) \tag{7a} \]

\[ \theta^{(2)} = \frac{-2}{2\sigma - 1} \left( \frac{1 - \sigma}{\sigma + (1 - \sigma) \cos(bz)} \right) \left[ \sqrt{2\sigma - 1} \arctan\left( \sqrt{2\sigma - 1} \tan \frac{bz}{2} \right) \right], \tag{7b} \]

where \( \sigma = \sqrt{p_c/p_0}, \ b = 2\sqrt{2}/\sigma^2 \) with \( p_c = \pi \) is the critical power for the soliton propagation.

Based on Eq. (5) and Eq. (7b), the phase shift is the function of both the propagation distance \( z \) and the power \( p_0 \). Worthy of note, the zero-order term of the phase shift in Eq. (5) is related to the size of the lead glass. This reveals the effect of the nonlocality on the phase shift of the SNSOS. In lead glass the nonlocality is essentially infinite [2] but cut-off by its boundary. Therefore it could be predicted that larger size glass should owe higher phase modulation sensitivity. Figure 1 demonstrates the phase shift of SNSOS in function of \( p_0/p_c \). We do not provide the numerical result in the case of \( w_0/R_0 = 1/300 \) because of the computer source available.

We carry out the large phase shift experiment in cylindrical lead glass with two diameters of 15 mm and 3 mm but the same length of 60 mm. The heavily-doped glass has a high absorption coefficient \( \alpha = 0.07 \) cm\(^{-1} \) and a high refractive index \( n_0 = 1.9 \). The other parameters are the same with those in Ref. [2]. The experimental arrangement is detailed in Fig. 2. A double frequency YAG laser (Verdi 12) with the wavelength of 532 nm is coupled into a Mach-Zehnder interferometer. The signal beam on one arm of the interferometer is focused by the lens \( L_1 \) onto the lead glass with beam width of 50 \( \mu \) m. The output soliton is imaged by \( L_2 \) onto CCD. The other arm contains a beam telescope, comprised by \( L_3 \) and \( L_4 \), adjusted to give a collimated, large diameter beam to act as a phase reference. Considering the difference of the refractive index between the lead glass and the air, a time delay is used in the reference optical path to compensate the optical length. The inset of Fig. 3 shows representative interference fringes with sharp contrast.

Along with the increase of the SNSOS power, the phase shift of the SNSOS increases considerably, while the phase of the reference beam keeps fixed. Therefore the interference fringes move observably. The stars in Fig. 3 designate the centers of a tracked fringe. The equidistant movement of the star indicates a linear modulation of the SNSOS phase by its power.

Considering the critical power is measured to be 260 mW, we change the input power from 190 mW to 340 mW by turning the tunable attenuator TA with power interval of 3.6 mW. Following the procedure in the treatment of the interference fringes in Fig. 3, we obtain the phase shift in function of the input power in lead glasses showed in Fig. 4. Since the linear fittings of the experimental data have slopes of 0.33 and 0.29, the \( \pi \) phase shifts are modulated by 9.5 mW power change in big glass bar and 10.5 mW power change in small glass bar respectively. In both case, the power changes are less...
than 5% of the soliton critical power and therefore the beams almost maintain the form of solitons when the power is slightly changed. Although the effect of the diameter of the glass on the phase shift is less than the theoretical predictions, the modulation sensitivity suggested by the experimental results is far higher than that predicted by the theoretical curves in Fig. 1. Segev et al. were in the similar situation when they numerically calculated the elliptic solitons [2] and measured the soliton steering driven by the boundary force [16]. The nonlinearity in their experiment is higher and more anisotropic than the calculated thermal response. The biggest steering data was three times than the theoretical prediction [16]. They presumed an additional mechanism in lead glass, the birefringence induced by thermal stress, giving rise to an increased $\Delta n$.

In conclusion we investigate the large phase shift of the SNSOS in lead glass. The experimental result verifies that an output phase shift of $\pi$ can be linearly modulated by a power change of about 10 mW, which is less than 5% of the soliton critical power. The effective producing of $\pi$ phase shift is significant to realize the treatment and control of the optical signal based on interference principle. Additionally the modulation sensitivity is higher in bigger size glass bar. This is the manifestation that the large phase shift of the SNSOS stems essentially from the strong nonlocality.

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