Goos-Hänchen Shifts of Partially Coherent Light Fields

Li-Gang Wang,[1,2,3,*] Shi-Yao Zhu,[4,†] and M. Suhail Zubairy[1,4,‡]
[1]Institute for Quantum Science and Engineering (IQSE) and Department of Physics and Astronomy, Texas A&M University, College Station, Texas 77843-4242, USA
[2]Department of Physics, Zhejiang University, Hangzhou 310027, China
[3]The National Center for Mathematics and Physics, KACST, P.O. Box 6086, Riyadh 11442, Saudi Arabia
[4]Beijing Computational Science Research Center, Beijing, 100084, China
(Received 26 June 2013; published 26 November 2013)

The Goos-Hänchen (GH) shift refers to a lateral displacement (from the path expected from geometrical optics) along an interface in totally internal reflection. This phenomenon results from a coherence effect. In order to bring to light the role of coherence, the reflection of partially coherent light fields was investigated within the framework of the theory of coherence. A formal expression for the GH shifts of partially coherent light fields is obtained in terms of Mercer's expansion. It is shown that both the spatial coherence and the beam width have an important effect on the GH shift, especially near the critical angles (such as totally reflection angle). These results are important to observe the GH shifts of the beams with imperfect coherence, like x-ray and matter-wave beams.

DOI: 10.1103/PhysRevLett.111.223901 PACS numbers: 42.60.Jf, 42.25.Gy, 42.25.Kb

When a bounded light beam is totally reflected from a planar interface, there is a lateral displacement from the path expected from geometrical optics. This is known as the longitudinal Goos-Hänchen (GH) shift [1]. This shift results from a coherence effect, and it is explained as [2] the different transverse wave vectors of a bounded light beam undergoing different phase changes, the sum of these waves forming a reflected beam with a lateral shift. Recently, the GH shift has also been seen as the sum of Renard's conventional energy flux plus a self-interference shift, which originates from the interference between the incident and reflected beams [3]. Furthermore, the classical Fresnel formulas for laws of refraction and reflection are discovered to be not applicable to partially coherent sources [4]. These investigations indicate that coherence should be very important to the GH shift.

However, recent investigations [5–9] have raised an important question: Does the spatial coherence influence the GH shift? The experiments, performed by Loffler *et al.* [6] and Merano *et al.* [7] did not demonstrate this effect, but the earlier experiment [10] observed the large difference between the measured GH shift of a partially coherent light-emitting diode light and the theoretical result of a coherent light, without a satisfactory explanation. To solve this issue is significant not only to the optical science, but also to other fields that involve the coherent wave phenomena, such as the GH shifts of neutrons [11,12], electrons [13,14], and spin waves [15]. The perfect coherent sources are hard to obtain for x-ray beams [16] and matter-wave beams [12,17]. As emphasized in Ref. [12], the observed GH shift may be used to accurately determine the coherence properties of the sources. Therefore, it is necessary to study this issue thoroughly and reveal the role of spatial coherence on the GH shift.

In this Letter, the scalar theory of coherence is employed to investigate the GH shift. First, a key formula to calculate the GH shift of partially coherent fields (PCFs) is derived in terms of the mode expansion. Then the physical mechanism, about the dependence of the GH shift on both the spatial coherence and beam width, is explained. Finally, an experimental proposal is suggested to show the impact of the spatial coherence on the practical GH shift near the critical angle.

A two-dimensional PCF is considered in the current issue. The cross-spectral density (CSD), $W(x_1, z_1; x_2, z_2; \nu)$, is employed to describe the propagation of the PCF, where (x_1, z_1) and (x_2, z_2) are the two points in the fields, and ν is the frequency of light. Here ν is omitted for simplicity. Then $W(x_1, z_1; x_2, z_2)$ can be expressed in the form of Mercer's expansion [18]

$$W(x_1, z_1; x_2, z_2) = \sum_m \beta_m \psi_m^*(x_1, z_1)\psi_m(x_2, z_2), \quad (1)$$

where ψ_m are the eigenfunctions and $\beta_m \geq 0$ are the eigenvalues. Equation (1) is also rewritten as

$$W(x_1, z_1; x_2, z_2) = \sum_n \beta_m W^{(m)}(x_1, z_1; x_2, z_2), \quad (2)$$

where $W^{(m)}(x_1, z_1; x_2, z_2) = \psi_m^*(x_1, z_1)\psi_m(x_2, z_2)$ represents the CSD of a coherent field. When PCFs are reflected at the interface ($z_{1,2} = z$) between two media, each mode ψ_m experiences a lateral shift, Δ_m. Therefore, the reflected CSD for ψ_m, at the interface, is given by

$$\begin{aligned} W_r^{(m)}(x_1, z_1; x_2, z_2) &= W_r^{(m)}(x_1, z; x_2, z) \\ &= |\bar{r}(\theta, \delta\theta_m)|^2 \psi_m^*(x_1 - \Delta_m, z) \\ &\quad \times \psi_m(x_2 - \Delta_m, z), \end{aligned} \quad (3)$$

where $\delta\theta_m$ is the angular spread of the mth mode, and $\bar{r}(\theta, \delta\theta_m)$ is the averaged reflection coefficient within $\delta\theta_m$ around the angle of incidence, θ. Since $\delta\theta_m$ may become very broad for a large m, the first-order Taylor expansion on the reflection coefficient r around θ is invalid [19]. Thus, Δ_m are different for different modes due to the size effect of each mode. Note that Δ_m are also different from the formula, $\Delta_{\mathrm{FT}} = -\mathrm{Re}[i(\partial \ln r/\partial\theta)]$ or $-(\lambda d\phi_r/2\pi d\theta)$, which is based on the stationary phase method under the approximation of the first-order Taylor expansion [2,20,21], here ϕ_r is the phase of r. Therefore, the total reflected CSD of a PCF is expressed as

$$W_r(x_1, z; x_2, z) = \sum_m w_m(\theta, \delta\theta_m)\psi_m^*(x_1 - \Delta_m, z) \times \psi_m(x_2 - \Delta_m, z), \quad (4)$$

where $w_m(\theta, \delta\theta_m) = \beta_m|\bar{r}(\theta, \delta\theta_m)|^2$ represents the weight of the mth reflected mode. Then the intensity of the reflected beam is

$$I_r(x, z) = \sum_m w_m(\theta, \delta\theta_m)|\psi_m(x - \Delta_m, z)|^2. \quad (5)$$

From the normalized first moment of a light field [22,23], $\Delta = \int xI_r(x, z)dx/\int I_r(x, z)dx$, the resultant GH shift is obtained as follows:

$$\Delta = \frac{\sum_m w_m(\theta, \delta\theta_m)\Delta_m}{\sum_m w_m(\theta, \delta\theta_m)}, \quad (6)$$

where the condition, $\int|\psi_n(x, z; \nu)|^2dx = 1$, is used. Equation (6) is a formal expression for calculating the GH shift of a PCF. Since the reflection coefficient r has not been specified to a particular system, Eq. (6) may be applied for calculating the lateral shifts of PCFs in any system [24], including both the partial and total reflections, for example, as discussed in Refs. [20,22,26]. This equation is different from that in Refs. [5,6]. In Refs. [5,6], all the lateral shifts Δ_m are assumed to be equal to Δ_{FT}, so that $\Delta = \Delta_{\mathrm{FT}}$. However, it is not true for PCFs, especially for the incoherent sources. In fact, it is shown (in the below discussion) that as m increases, there is a large difference between Δ_m and Δ_{FT}. Even for a coherent beam, Δ_m do change due to the finite-size effect of beams [25,27]. Therefore, the exact formulae for Δ_m for each mode is defined as [22,23,26]

$$\Delta_m = \int x|\psi_m^r(x, z)|^2dx \Big/ \int|\psi_m^r(x, z)|^2dx, \quad (7)$$

where ψ_m^r is the mth reflected field at the interface. For an incoherent light, all the contributions Δ_m to the resultant GH shift Δ must be included. Moreover, the weight factor, w_m, contains the value of β_m, which also strongly depends on the coherence properties of PCFs.

In order to demonstrate the lateral shift Δ_m of each mode of PCFs in Eq. (6), let us briefly review a famous example: a Gaussian shell-model (GSM) beam, which is an excellent model for describing PCFs [18]. The normalized eigenfunctions and eigenvalues of GSM beams are given by [18] (see also Refs. [28,29])

$$\psi_m(x) = \left(\frac{2c}{\pi}\right)^{1/4}\frac{1}{(2^m m!)^{1/2}}H_m[x(2c)^{1/2}]e^{-cx^2}, \quad (8)$$

and $\beta_m = A^2[\pi/(a + b + c)]^{1/2}[b/(a + b + c)]^m$, where $a = (4\sigma_s^2)^{-1}$, $b = (2\sigma_g^2)^{-1}$, $c = [a^2 + 2ab]^{1/2}$, and H_m are the Hermite polynomials. Here σ_s and σ_g are the beam half-width and spectral coherence width of PCFs, respectively. The ratio of β_m to β_0 is [18]

$$\beta_m/\beta_0 = [(q^2/2) + 1 + q[(q/2)^2 + 1]^{1/2}]^{-m}, \quad (9)$$

where $q = \sigma_g/\sigma_s$ is a measure of the degree of global coherence of a GSM source. Obviously, for $q \gg 1$, $\beta_m/\beta_0 \approx q^{-2m} \ll 1$ for $m > 0$. Hence, the beam can be well approximated by the lowest-order mode. However, for $q \ll 1$, $\beta_m/\beta_0 \approx 1 - mq$. Thus, for an incoherent light, a large number of modes (of the order $1/q$) are needed to represent the light field adequately.

Since each mode is perfectly coherent, its shift can be obtained from the coherent angular-spectral theory [20,22,26] under a certain angle of incidence, as illustrated in Fig. 1(a). From Eq. (8), its angular spectrum, $\tilde{\psi}_m(k_x)$, can be obtained via a Fourier transformation. For an inclined incidence, $\tilde{\psi}_m(k_x)$ becomes $\tilde{\psi}_m(k_x - k_{x0})$ with the replacements $\sigma_s \to \sigma_s\sec\theta$ and $\sigma_g \to \sigma_g\sec\theta$, where k_x is the transverse component of the wave vector $\mathbf{k}$ in the prism, and $k_{x0} = k\sin\theta$. Therefore, ψ_m^r is given by

$$\psi_m^r(x) = \frac{1}{\sqrt{2\pi}}\int r(k_x)\tilde{\psi}_m(k_x - k_{x0})\exp[ik_x x]dk_x. \quad (10)$$

Then, from Eq. (7), all shifts Δ_m can be obtained. In the following calculations, the refractive index of the prism is $n = 1.514$ at wavelength $\lambda = 675$ nm, and the critical angle of the totally internal reflection is $\theta_c = 41.34°$. Here, only the result for the transverse magnetic (TM) polarization is presented due to the similarity between TM and transverse electric (TE) cases [30].

Effect of spatial coherence.—Figures 1(b) and 1(c) show the typical dependence of the lateral shifts Δ_m on the spatial coherence (q) under different cases: (b) $\theta = 41.5°$ and (c) $\theta = 45°$. Here, $\sigma_s = 0.1$ mm ($\gg \lambda$) in these two cases. From Figs. 1(b) and 1(c), the lateral shifts Δ_m near θ_c are strongly dependent on q. For $m = 0$, the value Δ_0 slightly increases when q is gradually close to 0.1, but it decreases as q further decreases. For a large m, the changes of Δ_m become more dramatic with the decreasing of q, and more oscillations appear due to the fact that the part components of $\tilde{\psi}_m(k_x - k_{x0})$ have been cut off below θ_c as a result of the broadening angular spectrum of $\tilde{\psi}_m(k_x - k_{x0})$ with the decreasing of q. When θ is far above θ_c, see Fig. 1(c), the values Δ_m have a greater change for the large value of m. Thus, there exists a difference between the coherent

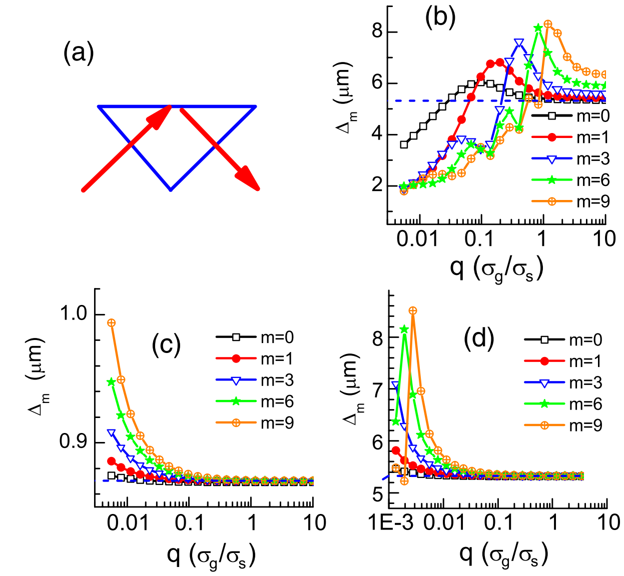

FIG. 1 (color online). (a) The schematic of total reflection from a prism. (b)–(d) The dependence of the lateral shifts Δ_m on the spatial coherence (q) at different θ: $\theta = 41.5°$ (b),(d) and $\theta = 45°$ (c). The blue dashed lines in (b),(c),(d) denote the values Δ_{FT}. In (b),(c) $\sigma_s = 0.1$ mm, and in (d) $\sigma_s = 2$ mm.

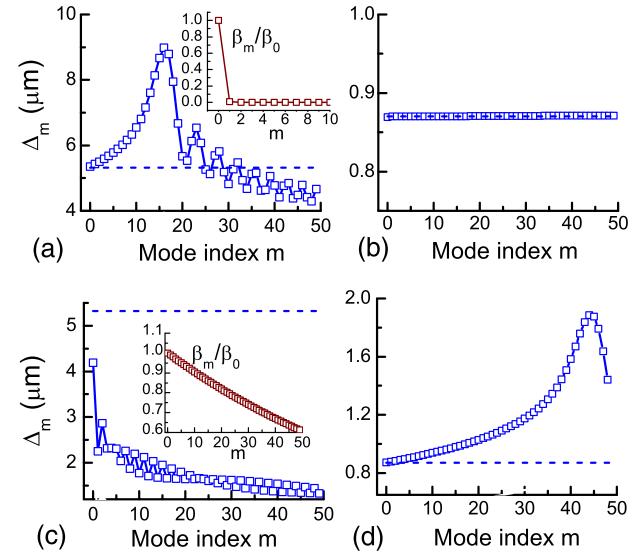

FIG. 2 (color online). The lateral shifts Δ_m as a function of m under two limits: $q = 10$ (a),(b) and $q = 0.01$ (c),(d), at $\theta = 41.5°$ (a),(c) and $\theta = 45°$ (b),(d). Insets in (a),(c) show the value of β_m/β_0 as a function of m for $q = 10$ and $q = 0.01$, respectively. The dashed lines in (a)–(d) denote the values Δ_{FT}.

and incoherent limits [8,31]. By comparing Fig. 1(b) with Fig. 1(c), it is further found that the changes of Δ_m near θ_c are more remarkable than for θ being far above θ_c.

Figure 1(d) plots another situation for the dependence of Δ_m on q when $\theta = 41.5°$, but $\sigma_s = 2$ mm. Although θ is near to θ_c, the changes of Δ_m are considerably small as long as $q > 0.01$. This is due to the suppressing effect of beam width ($2\sigma_s$) on Δ_m discussed below. In Fig. 1(d), it is also shown that there is still a large difference between Δ_m and Δ_{FT} in the incoherent limit ($q < 0.01$). As m increases, some oscillations still appear for a sufficient small q.

Figure 2 further shows the changes of Δ_m as a function of m under the two limits: $q = 10$ (coherent) and $q = 0.01$ (incoherent). For the fully coherent limit ($q \gg 1$), when θ is close to θ_c [see Fig. 2(a)], Δ_m vary dramatically as m increases, while when θ is far above θ_c [see Fig. 2(b)], Δ_m are nearly independent of m and they are overlapped with the corresponding value of Δ_{FT}. Thus, in the full-coherent limit, Δ_m are independent of m only when θ is far away from θ_c. Meanwhile, only the lateral shift Δ_0 of the lowest mode ($m = 0$) mainly contributes to the resultant shift Δ, because β_m decrease quickly for $m > 0$, see the inset in Fig. 2(a). For the completely incoherent limit ($q \ll 1$), see Figs. 2(c) and 2(d), whether θ is close to or far away from θ_c, Δ_m do vary as m increases, and the contributions of the higher-order modes must be included since β_m change very slowly for $m > 0$, see the inset in Fig. 2(c). This leads to the resultant shift deviated from the full-coherent limit.

Effect of beam width ($2\sigma_s$).—The beam width of PCFs also plays a role on the GH shift, since the effective width ($2\sigma_m^{\mathrm{eff}}$) of every ψ_m is related to both σ_s and σ_g [18]. From Eq. (8), one can obtain $\sigma_m^{\mathrm{eff}} = \sqrt{2m+1}\sigma_s/[1 + (4/q^2)]^{1/4}$ and $\delta\theta_m = (180\sqrt{2m+1}/\pi k \sigma_s)[1 + (4/q^2)]^{1/4}$ (in the unit of degree). For a fixed q, if σ_s increases, then σ_m^{eff} increases but $\delta\theta_m$ decreases. It means that increasing σ_s suppresses the effect of q on the shift Δ_m. By comparing Figs. 1(b) and 1(d), it is seen that increasing σ_s leads to weaken the effect of spatial coherence on the GH shift. In principle, an incoherent field has an infinite number of modes; for any q, as $m \to \infty$, the shift Δ_∞ should be zero because of $\sigma_m^{\mathrm{eff}} \to \infty$ and $\delta\theta_m \to \infty$. It should be emphasized that, for a coherent beam, the effect of beam width has been investigated in the early literature [19,25] and has been demonstrated experimentally [32]. Therefore, it is expected that the beam width also has an effect on the GH shift for PCFs.

Figure 3 shows the detailed effect of σ_s on Δ_m near θ_c. It shows that, see Fig. 3(a), even for the full-coherent limit ($q = 10$), when σ_s is small enough (< 0.3 mm but $\gg \lambda$), the values Δ_m begin to be significantly different from Δ_{FT}, and the difference becomes larger as m increases. Remember that it is only the lowest mode that dominates the shift Δ in the full-coherent limit, because $\beta_m \to 0$ for $m > 0$. However, in the incoherent limit ($q = 0.01$), see Fig. 3(b), when σ_s is larger than 2 mm, the difference between Δ_m and Δ_{FT} gradually disappears due to the suppressing effect of σ_s on Δ_m, while for $\sigma_s < 2$ mm, Δ_m change dramatically and are very different from Δ_{FT}.

In fact, the role of σ_s on Δ_m for a small q is similar to the role of q on Δ_m for a small σ_s, see Figs. 1(b) and 3(b). On comparing Fig. 3(a) with Fig. 1(d), it further shows that the role of σ_s on Δ_m for a large q is similar to the role of q on

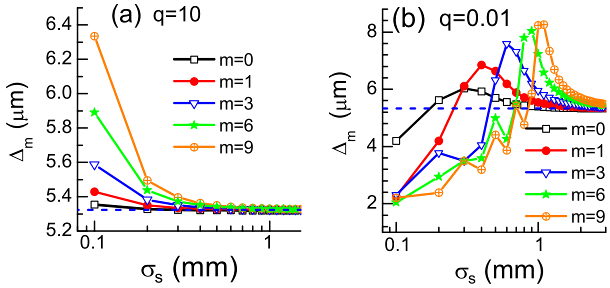

FIG. 3 (color online). Effect of σ_s on the shifts Δ_m of each mode under $q = 10$ (a) and $q = 0.01$ (b), with $\theta = 41.5°$.

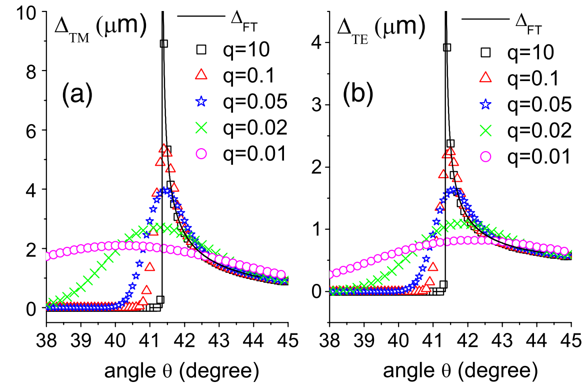

FIG. 4 (color online). The resultant GH shifts, for (a) TM cases and (b) TE cases, as a function of the angle of incidence under different q with a fixed $\sigma_s = 0.2$ mm. Here $\theta_c = 41.34°$.

Δ_m for a large σ_s. Therefore, both σ_s and q have the equivalent role on the GH shift.

Now the roles of σ_s and q on the GH shift have been known, and how or why they affect the shift Δ has been explained. However, it is inconvenient for using Eq. (6) to obtain Δ since it is time consuming to calculate all shifts Δ_m when q is very small. For example, if $q = 0.01$, it needs 100 modes at least. There is a much more realistic method for directly obtaining Δ. Based on the previous investigation, the intensity expression of the reflected PCFs, at the interface ($z = 0$), is given by [31],

$$\begin{aligned} I_r(x,0) &= W_r(x,0;x,0) \\ &= \frac{1}{2\pi}\iint r^*(k_{x1})r(k_{x2})W_i(k_{x1},0;k_{x2},0) \\ &\quad \times \exp[-i(k_{x1}-k_{x2})x]dk_{x1}dk_{x2}, \end{aligned} \tag{11}$$

where $W_i(k_{x1}, 0; k_{x2}, 0)$ is the incident CSD in the spatial angular-frequency domain at $z = 0$. Substituting Eq. (11) into the definition of Δ, we can obtain the GH shift of PCFs by the numerical method.

Finally, let us briefly discuss how to experimentally demonstrate the effect of spatial coherence on the GH shift, since the experiments [6,7] have not revealed this effect. In the experiments [6,7], the authors increased σ_s for obtaining a small q, which in fact suppresses the effect of q on the GH shift. From the above discussion, it is known that the large σ_s weakens the effect, and near θ_c, the spatial coherence has a larger effect. Thus, choosing a small and fixed σ_s is beneficial for observing the effect near θ_c in the experiment. Figure 4 presents the dependence of the absolute GH shifts, for both TM and TE cases, on its spatial coherence for experimental reference. Here $\sigma_s = 0.2$ mm, and $q = 10, 0.1, 0.05, 0.02$, and 0.01. From Fig. 4, for a full coherent light ($q = 10$), there are nonzero GH shifts above θ_c but zeros below θ_c, and the shifts Δ in this situation, for both TM and TE cases, are overlapped with the corresponding curves of $\Delta_{\rm FT}$, respectively. However, for a PCF or an incoherent light field, the shifts above θ_c may be smaller or larger than $\Delta_{\rm FT}$. More interestingly, the lateral shifts below θ_c, for both TM and TE cases, are no longer equal to zero. This is a distinct result, for PCFs, which is completely different from the full-coherent prediction. In fact, this effect has been observed in a recent experiment [10], where a nonzero lateral shift below θ_c is measured without an appropriate explanation. The nonzero GH shifts of PCFs below θ_c are very similar (qualitatively not quantitatively) to the effect of the narrow beam width on the GH shifts [25,33]. It should also be pointed out that the profiles of the reflected incoherent fields are unchanged even if θ is close to (or below) θ_c. Since the curves in Fig. 4 have the same characteristic as with other experiments [23,34,35], we hope this suggestion could lead to a direct experimental observation in the system of the total internal reflection.

In summary, the formal expression (6) of the GH shift of PCFs, which was obtained by using the exact theory of coherence, reveals its dependence on both the spatial coherence and beam width. This result can explain why the recent experimental results did not show this effect [6,7]. Meanwhile, a potential experiment is discussed for demonstrating this effect and displaying a distinct effect for experimental verification. These effects are very important to the applications of the GH shift in nano- and micro-scaled structures [34,36,37], where the light sources are usually focused into the small region (about a few wavelengths or less) and coherence plays a significant role in the shift. The results also have important impact for other fields, where the fully coherent sources are usually not available. For example, the x-ray beams are usually partially coherent [38], the results presented here may truly explain a discrepancy of the GH shift of the x ray near the critical angle in its Bragg diffraction [16]. Furthermore, the result indicates that, in the future, observation of the GH shift of matter-wave beams, such as neutron beams [11,12], atom beams [17], and electron beams [13,14], should include the impact of the spatial coherence of the sources, especially for highly precise measurements.

This research is supported by NPRP Grant No. 4-346-1-061 by the Qatar National Research Fund (QNRF) and a grant from King Abdulaziz City for Science and Technology (KACST). This work is also supported by

NSFC Grants (No. 61078021, No. 11174026, and No. 11274275), and by the National Basic Research Program of China (Grants No. 2012CB921602,3 and No. 2011CB922203).

*sxwlg@yahoo.com
†syzhu@csrc.ac.cn
‡zubairy@physics.tamu.edu

[1] F. Goos and H. Hänchen, Ann. Phys. (Berlin) **436**, 333 (1947).
[2] K. Artmann, Ann. Phys. (Berlin) **437**, 87 (1948).
[3] X. Chen, X.-J. Lu, P.-L. Zhao, and Q.-B. Zhu, Opt. Lett. **37**, 1526 (2012).
[4] M. Lahiri and E. Wolf, Phys. Rev. A **86**, 043815 (2012).
[5] A. Aiello and J. P. Woerdman, Opt. Lett. **36**, 3151 (2011).
[6] W. Löffler, A. Aiello, and J. P. Woerdman, Phys. Rev. Lett. **109**, 213901 (2012).
[7] M. Merano, G. Umbriaco, and G. Mistura, Phys. Rev. A **86**, 033842 (2012).
[8] L. G. Wang and K. H. Liu, Opt. Lett. **37**, 1056 (2012).
[9] L. G. Wang, S. Y Zhu, and M. S. Zubairy, arXiv:1304.5585.
[10] H. G. L. Schwefel, W. Köhler, Z. H. Lu, J. Fan, and L. J. Wang, Opt. Lett. **33**, 794 (2008).
[11] V. K. Ignatovich, Phys. Lett. A **322**, 36 (2004).
[12] V.-O. de Haan, J. Plomp, T. M. Rekveldt, W. H. Kraan, A. A. van Well, R. M. Dalgliesh, and S. Langridge, Phys. Rev. Lett. **104**, 010401 (2010).
[13] C. W. J. Beenakker, R. A. Sepkhanov, A. R. Akhmerov, and J. Tworzydlo, Phys. Rev. Lett. **102**, 146804 (2009).
[14] Z. Wu, F. Zhai, F. M. Peeters, H. Q. Xu, and K. Chang, Phys. Rev. Lett. **106**, 176802 (2011).
[15] Y. S. Dadoenkova, N. N. Dadoenkova, I. L. Lyubchanskii, M. L. Sokolovskyy, J. W. Klos, J. Romero-Vivas, and M. Krawczyk, Appl. Phys. Lett. **101**, 042404 (2012).
[16] K. Tamasaku and T. Ishikawa, Acta Crystallogr. Sect. A **58**, 408 (2002).
[17] J. H. Huang, Z. L. Duan, H. Y. Ling, and W. P. Zhang, Phys. Rev. A **77**, 063608 (2008).
[18] L. Mandel and E. Wolf, *Optical Coherence and Quantum Optics* (Cambridge University Press, Cambridge, England, 1995).
[19] B. R. Horowitz and T. Tamir, J. Opt. Soc. Am. **61**, 586 (1971).
[20] C. F. Li, Phys. Rev. Lett. **91**, 133903 (2003).
[21] R. Simon and T. Tamir, J. Opt. Soc. Am. A **6**, 18 (1989).
[22] I. V. Shadrivov, A. A Zharov, and Y. S. Kivshar, Appl. Phys. Lett. **83**, 2713 (2003).
[23] F. Bretenaker, A. Le Floch, and L. Dutriaux, Phys. Rev. Lett. **68**, 931 (1992).
[24] When θ is larger than the critical angle θ_c in the total reflection, the lateral shift is called the GH shift. When θ is less than θ_c, the lateral shift is zero from the prediction of the stationary phase method, but it is nonzero only for a very narrow coherent beam (a few wavelengths), due to the fact that the part of the angular spectrum is above the critical angle [25]. For systems of partial reflection, like a slab system [20,26], it's better to be called the lateral shift (or the spatial beam shift).
[25] H. M. Lai, F. C. Cheng, and W. K. Tang, J. Opt. Soc. Am. A **3**, 550 (1986).
[26] F. Wang and A. Lakhtakia, Opt. Commun. **235**, 107 (2004).
[27] D. Golla and S. D. Gupta, Pramana J. Phys. **76**, 603 (2011).
[28] F. Gori, Opt. Commun. **34**, 301 (1980).
[29] J. T. Foley and M. S. Zubairy, Opt. Commun. **26**, 297 (1978).
[30] Except for the magnitudes, the behaviors of the lateral shifts Δ_m, for both the TM and the TE cases, are similar to each other, in the system of total internal reflection.
[31] L.-Q. Wang, L.-G. Wang, S.-Y. Zhu, and M. S. Zubairy, J. Phys. B **41**, 055401 (2008).
[32] J. J. Cowan and B. Aničin, J. Opt. Soc. Am. **67**, 1307 (1977).
[33] J. F. Gray and A. Puri, Phys. Rev. A **75**, 063826 (2007).
[34] X. Yin, L. Hesselink, Z. Liu, N. Fang, and X. Zhang, Appl. Phys. Lett. **85**, 372 (2004).
[35] J. Unterhinninghofen, U. Kuhl, J. Wiersig, H.-J. Stöckmann, and M. Hentschel, New J. Phys. **13**, 023013 (2011).
[36] I. V. Soboleva, V. V. Moskalenko, and A. A. Fedyanin, Phys. Rev. Lett. **108**, 123901 (2012).
[37] A. Lakhtakia, Proc. SPIE Int. Soc. Opt. Eng. **5218**, 29 (2003).
[38] K. A. Nugent, Adv. Phys. **59**, 1 (2010).