ABSTRACT

We present both theoretical and experimental study of longitudinal optical binding between two nanopaticles or microparticles placed in two counter-propagating Bessel beams. We consider coherent linearly polarized beams and we rotate polarization of one beam so that both beams interfere or do not interfere. We present an experimental demonstration of the remarkable changes in the behaviour of both spheres leading to freezing of the structure if both beams interfere.

Keywords: optical binding, optical forces, optical tweezers, optical trapping, Bessel beams

1. INTRODUCTION

Optical binding has been presented about 20 years ago by Burns et al.\textsuperscript{1} and it denotes an interesting type of the light-matter interaction for mesoscopic objects which leads to their self-arrangement into so called “optical matter”.\textsuperscript{2} This intriguing self-arrangement is based upon delicate equilibrium between the optical forces resulting both from the incident beam and from the light re-scattered by the objects which then results in well defined stable equilibrium positions for the particles concerned. The mechanism is different from the well-known optical trapping because here no predefined potential wells (optical traps) exist before the objects are placed to the field. The presence of the objects modifies the field distribution so that possible optical traps are created for interacting objects. The aim in the field of optical binding is to understand these phenomena in depth to pave the way for controlled self-assembly of colloidal particles solely by light.

In this paper we concentrate on longitudinal optical binding\textsuperscript{3,4} occurring along the beam propagation while laterally the objects are confined by the mechanism of optical trapping in the high intensity core of the beam. In contrast to the lateral (transverse) binding studied by Burns et al.,\textsuperscript{1} the longitudinal binding forms optical matter in three dimensions far from any surface and has the potential for long-range particle arrangement. Such longitudinal optical binding has been studied mainly in Gaussian beams.\textsuperscript{5,6} In our previous studies we have already demonstrated long-range 1-D optical binding of micrometer and sub-micrometer sized objects placed within propagation invariant (non-diffracting) Bessel light beams\textsuperscript{7,8} in a counter-propagating geometry. In such configuration the optical forces corresponding to light scattering by the objects are dominant over gradient forces because the Bessel beam has negligible intensity variation along its propagation axis. The beam is not significantly disturbed by the inserted objects due to its self-reconstructing property\textsuperscript{9,10} leading to an increased number of objects being organized in the beam.\textsuperscript{11–13}

In this paper we focus on the optical binding if both beams interfere strongly or almost do not interfere. We present a simplified theory of optical binding of two nanoparticles placed into standing wave and compare the results with the configuration where the beams do not interfere. Due to the existing experimental setup we studied such phenomenon in Bessel beams theoretically and also experimentally.

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2. OPTICAL BINDING OF NANOPARTICLES

2.1 The evaluation of field scattered by two dipoles

A nanoparticle can be understood as an induced dipole or so called Rayleigh particle. Its polarizability \( \alpha \) gives a measure of induced electric dipole momenta \( \mathbf{p} \) through linear relationship \( \mathbf{p} = \alpha \mathbf{E} \) with \( \mathbf{E} \) being the electric field incident on the dipole. Let us assume two identical induced dipoles (approximated as tiny spheres of radius \( a \)) with scalar polarizability \( \alpha \) with included radiative reaction term: \(^{14}\)

\[
\alpha = \frac{\alpha_0}{1 - \frac{3}{4} k a^3 \frac{\alpha_0}{4 \pi \varepsilon}} \simeq \alpha_0 \left( 1 + \frac{k^3}{6 \pi \varepsilon} \alpha_0 \right) = \alpha' + i \alpha'',
\]

where \( \alpha' \) or \( \alpha'' \) denotes the real or imaginary part of \( \alpha \), respectively, and \( \alpha_0 \) comes from the Lorentz-Lorenz equation: \(^{15}\)

\[
\alpha_0 = 4 \pi a^3 \frac{m^2 - 1}{m^2 + 2},
\]

where \( m = \eta_p/\eta_m \), \( \varepsilon = \varepsilon_0 \eta_m^2 \) is permittivity of medium, \( \eta_p \) and \( \eta_m \) are refractive indices of particles and surrounding medium, respectively, and \( \varepsilon_0 \) is the permittivity of vacuum.

Let us focus on the interaction between two dipoles which are located on a common axis of two counter-propagating beams as depicted in Fig. 1. We assume the dipoles are located on the optical axis and forces between the dipoles are evaluated only in this direction. The induced dipoles of both particles are

\[
\mathbf{p}(\mathbf{r}_A) = \alpha \mathbf{E}(\mathbf{r}_A), \quad \mathbf{p}(\mathbf{r}_B) = \alpha \mathbf{E}(\mathbf{r}_B).
\]

Let us simplify this notation using \( \delta \mathbf{E} \equiv \mathbf{E}(\mathbf{r}_A) \) (resp. \( \delta \mathbf{E} \equiv \mathbf{E}(\mathbf{r}_B) \)) and \( \mathbf{p}(\mathbf{r}_A) \equiv \delta \mathbf{p} \) (resp. \( \mathbf{p}(\mathbf{r}_B) \equiv \delta \mathbf{p} \)) to write:

\[
\delta \mathbf{p} = \alpha \delta \mathbf{E}, \quad \delta \mathbf{p} = \alpha \delta \mathbf{E}.
\]

The electric field \( \delta \mathbf{E}' \) at \( \mathbf{r}_A \) caused by the radiating dipole \( \mathbf{B} \) located at \( \mathbf{r}_B \) (having the dipole moment \( \delta \mathbf{p} \)), is given by \( \delta \mathbf{E}' = \mathbf{G}(\mathbf{r}_A - \mathbf{r}_B) \delta \mathbf{p} \), where the field propagator \( \mathbf{G}(\mathbf{r}) \) acting on dipole momentum of the particle is given as: \(^{15}\)

\[
\mathbf{G}(\mathbf{r}) = \frac{\exp(i k \mathbf{r})}{4 \pi \varepsilon r^3} \left[ (-k^2 r^2 - 3 i k r + 3) \frac{\mathbf{r} \otimes \mathbf{r}}{r^2} + (k^2 r^2 + i k r - 1) \mathbf{1}_3 \right],
\]

where \( k = 2 \pi / \lambda \), \( \otimes \) represents tensor product of two vectors and \( \mathbf{1}_3 \) is identity matrix. The total exciting field given by an incident field and the radiation of two dipoles is a solution of a set of equations (4) and the following ones:

\[
\delta \mathbf{E} = \delta \mathbf{E}^{\text{inc}} + \mathbf{G}(\mathbf{r}_A - \mathbf{r}_B) \delta \mathbf{p},
\]

\[
\delta \mathbf{E} = \delta \mathbf{E}^{\text{inc}} + \mathbf{G}(\mathbf{r}_B - \mathbf{r}_A) \delta \mathbf{p}.
\]

We assume that both dipoles \( \mathbf{A} \) and \( \mathbf{B} \) are located on \( z \) axis and their mutual distance is \( d \) as depicted in Fig. 1. Therefore, the non-zero components of the propagator \( \mathbf{G}(\mathbf{r}) \) with \( \mathbf{r} = (0, 0, d) \) are placed only on the diagonal (labelled similarly as in Ref.\(^{16}\)):

\[
\mu(d) \equiv \mathbf{G}_{xx}(0, 0, d) = \mathbf{G}_{yy}(0, 0, d) = \frac{\exp(i k d)}{4 \pi \varepsilon d^3} [k^2 d^2 + i k d - 1],
\]

\[
\lambda(d) \equiv \mathbf{G}_{zz}(0, 0, d) = \frac{\exp(i k d)}{4 \pi \varepsilon d^3} [-2 i k d + 2].
\]
Substitution of these terms into equation (7) gives the total exciting field at the positions of the dipoles \( A \) and \( B \) as:

\[
\begin{align*}
A_{E_x} &= \frac{A_{E_{x \text{inc}}} + \mu\alpha B_{E_{x \text{inc}}}}{1 - \mu^2\alpha^2}, \\
B_{E_x} &= \frac{B_{E_{x \text{inc}}} + \mu\alpha A_{E_{x \text{inc}}}}{1 - \mu^2\alpha^2}, \\
A_{E_y} &= \frac{A_{E_{y \text{inc}}} + \mu\alpha B_{E_{y \text{inc}}}}{1 - \mu^2\alpha^2}, \\
B_{E_y} &= \frac{B_{E_{y \text{inc}}} + \mu\alpha A_{E_{y \text{inc}}}}{1 - \mu^2\alpha^2}, \\
A_{E_z} &= \frac{A_{E_{z \text{inc}}} + \lambda\alpha B_{E_{z \text{inc}}}}{1 - \lambda^2\alpha^2}, \\
B_{E_z} &= \frac{B_{E_{z \text{inc}}} + \lambda\alpha A_{E_{z \text{inc}}}}{1 - \lambda^2\alpha^2}.
\end{align*}
\] (9) (10) (11)

Because of the diagonal shape of \( \mathbf{G} \), the \( x, y, z \) components of total exciting fields are only composed from corresponding \( x, y, z \) components of incident fields. For the case of simplicity we will consider paraxial beams in the rest of the paper and therefore we will not consider fields with axial component \((A_{E_z} = B_{E_z} = 0)\).

### 2.2 The simplified solution for fields and forces

Depasse and Vigoureux\(^\text{16}\) presented a complete and also approximated solution of Eqs. (9-11) for plane wave and applied them to get acting forces without any approximation. They considered configuration with two dipoles having their connecting line perpendicular to the direction of incident plane wave. Here we consider a different configuration where the connecting line is parallel to the direction of two counter-propagating fields. The main distinction in our approach is application of physically justified approximations to the system of equations prior to the solution. This approach greatly simplifies the derivation and reveals the physical significance of individual force terms. The term \( \mu\alpha \) in Eq. (9) describes the field propagation (radiation) from one dipole (given by \( \alpha \)) to another. Its second power i.e. \( \mu^2\alpha^2 \) describes the contribution to the total exciting field at the position of one dipole caused by the radiation coming back from this dipole through the second dipole. This process corresponds to scattering of the second order and its effect can be usually neglected. Even higher powers of the term \( \mu\alpha \) occur in the equations dealing with forces in the longitudinal binding setup. We estimate its value assuming Rayleigh spherical particles with their centres separated by \( d \):

\[
\mu\alpha = C \exp(\text{i}kd) \left[ k^2 \frac{a^3}{d} + \text{i}k \frac{a^3}{d^2} - \frac{a^3}{d^3} \right], \text{ where } C = \frac{m^2 - 1}{m^2 + 2}.
\] (12)

For the commonly used dielectric particles \((n_p = 1.6)\) placed in water \((n_m = 1.332)\) the value of the factor \( C \) depending only on refractive indices is approximately \( C = 0.2 \). For Rayleigh particles and non-contact condition \((d > 2a)\) the magnitudes of all three terms in Eq. (12) differ: The magnitude of the first term, proportional to \( k^2 \), is smaller than 0.015, the second term, proportional to \( k \), is smaller than 0.026, and the third one is smaller than 0.041. Although the third term gives the highest contribution if the spheres are in contact \(|\mu\alpha|^2 \leq 0.02\), its magnitude decreases very rapidly \((d^{-3})\). Therefore the first term proportional to \( k^2 \) becomes dominant for distances bigger than \( \lambda/(2\pi) \) and results in \(|\mu\alpha|^2 \leq 10^{-4}\).

Therefore, the denominators in equations (9) and (10) for total exciting field are equal to unity and only the fields incident on each dipole occur on the right hand side of these equations. This attitude resembles the first Born approximation.\(^\text{17}\) Let us assume that both incident beams are linearly polarized in \( x \) or \( y \) direction and so we get from equations (9):

\[
\begin{align*}
A_{E_x} &= A_{E_{x \text{inc}}} + \mu\alpha B_{E_{x \text{inc}}}, \\
B_{E_x} &= B_{E_{x \text{inc}}} + \mu\alpha A_{E_{x \text{inc}}}, \\
A_{E_y} &= A_{E_{y \text{inc}}} + \mu\alpha B_{E_{y \text{inc}}}, \\
B_{E_y} &= B_{E_{y \text{inc}}} + \mu\alpha A_{E_{y \text{inc}}},
\end{align*}
\] (13) (14)

The general expression for a time-averaged force acting in \( z \) direction on an induced dipole has the following form:\(^\text{18}\)

\[
\langle F_z \rangle = \frac{1}{2} \Re \left\{ p_x^* \frac{\partial E_x}{\partial z} + p_y^* \frac{\partial E_y}{\partial z} + p_z^* \frac{\partial E_z}{\partial z} \right\}.
\] (15)
To obtain the optical forces acting on both considered interacting dipoles we need to express the spatial derivatives of these approximated total exciting fields (14) up to the first order of $\mu$:

$$\frac{\partial^2 E_x}{\partial z_A^2} = \frac{\partial^2 E_{x,\text{inc}}}{\partial z_A^2} + \frac{\partial \mu}{\partial z_A} \alpha E_{x,\text{inc}}, \quad \frac{\partial^2 E_x}{\partial z_B^2} = \frac{\partial^2 E_{x,\text{inc}}}{\partial z_B^2} + \frac{\partial \mu}{\partial z_B} \alpha E_{x,\text{inc}},$$  \tag{16}

$$\frac{\partial^2 E_y}{\partial z_A^2} = \frac{\partial^2 E_{y,\text{inc}}}{\partial z_A^2} + \frac{\partial \mu}{\partial z_A} \alpha E_{y,\text{inc}}, \quad \frac{\partial^2 E_y}{\partial z_B^2} = \frac{\partial^2 E_{y,\text{inc}}}{\partial z_B^2} + \frac{\partial \mu}{\partial z_B} \alpha E_{y,\text{inc}}.$$  \tag{17}

Putting the approximative fields and their spatial derivatives into (15), the following expressions of the resulting forces acting on induced dipoles $A$ and $B$ are obtained:

$$\langle F^A \rangle = \frac{1}{2} \Re \left\{ \left( \alpha A_{E_{j,\text{inc}}}^* \right) \frac{\partial A_{E_{j,\text{inc}}}}{\partial z_A} + \left( \alpha B_{E_{j,\text{inc}}} \right) \frac{\partial B_{E_{j,\text{inc}}}}{\partial z_A} \right\},$$  \tag{18a}

$$\langle F^B \rangle = \frac{1}{2} \Re \left\{ \left( \alpha B_{E_{j,\text{inc}}}^* \right) \frac{\partial B_{E_{j,\text{inc}}}}{\partial z_B} + \left( \alpha A_{E_{j,\text{inc}}} \right) \frac{\partial A_{E_{j,\text{inc}}}}{\partial z_B} \right\}.$$  \tag{18b}

where we accepted Einstein summation rule over the same indices $j$ denoting $x$ and $y$ components of the incident fields. Let us explain the individual parts of the force acting on the second induced dipole $B$. The first term is the most dominant because the radiation factor $\mu$ is absent here. It is equivalent to the force acting on a dipole $B$ induced by the incident field at $\mathbf{n}_k$ and interacting with this field. No contribution from the second dipole $A$ is included here. The second force term describes an induced dipole created by the field scattered from the dipole $A$ and placed to the incident field at $\mathbf{n}_B$. The third term corresponds to the situation when the dipole is induced by the incident field at $\mathbf{n}_B$ but interacting with the field scattered from dipole $A$. The last fourth term describes how the dipole at $B$ induced by the scattered field from the dipole $A$ interacts with the scattered field from dipole $A$. So the magnitude of this last term is roughly proportional to $|\mu| a^2$ and it may be neglected because in our case it is $10^{-2}$ times smaller (see expression (12)) than the second and the third force terms.

### 2.3 Longitudinal binding of nanoparticles in counter-propagating interfering Bessel beams

We will assume that the beam propagating along positive direction of $z$ axis is polarized along $x$ axis, and the polarization of the wave propagating against the direction of $z$ axis is tilted by $\phi$ with respect to the $x$ axis. Therefore if $\phi = \pi/2$, both waves do not interfere. For the purposes of this paper we will further assume that amplitude of the incident fields does not vary along the propagation axis. Since the nanoparticles are placed on the optical axis the incident Bessel beams, they can be approximated as plane waves of amplitudes $A_{E_{0+x}}^{\text{inc}}$ and $A_{E_{0-z}}^{\text{inc}}$ propagating with wavevector $k_z$ and $-k_z$, respectively. Therefore for particle placed at $z = z_A$ we obtain:

$$A_{E_{x,\text{inc}}} = A_{E_{0+x}}^{\text{inc}} \exp(ik_z z_A) + A_{E_{0-z}}^{\text{inc}} \cos \phi \exp(-i k_z z_A),$$  \tag{19}

$$A_{E_{y,\text{inc}}} = A_{E_{0+z}}^{\text{inc}} \sin \phi \exp(-i k_z z_A).$$  \tag{20}

Further we will assume identical beams and we set $A_{E_{0+x}}^{\text{inc}} = A_{E_{0-z}}^{\text{inc}} = A_{E_{0x}}^{\text{inc}}$. Thus the radiation pressures coming from both beams are compensated. Moreover, the nanoparticles are placed on the optical axis and therefore the radial profile of the beam is not important for the axial behaviour of the particles. Substituting these equations

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into (18a) and (18b) we obtain
\[
\langle AF \rangle = -\alpha' |E_0^{inc}|^2 \left\{ k_z \cos \phi \sin(2k_z z_A) - \frac{\alpha' k^2}{4\pi \varepsilon d} \left[ k_z \cos(kd) \sin(k_z d) + k \cos(k_z d) \sin(kd) \right] \right\},
\]
\[
\langle BF \rangle = -\alpha |E_0^{inc}|^2 \left\{ -k_z \cos(kd) \sin(k_z z_A + k_z z_B) + k \sin(kd) \cos(k_z z_A + k_z z_B) \right\},
\]
where the distance between the particles \( d = z_B - z_A \) and only the far field terms are kept. The other terms proportional to \( 1/d^2, 1/d^3, 1/d^4 \) were omitted here.

If \( \phi = \pi/2 \), the counter-propagating waves do not interfere and pure longitudinal optical binding occurs
\[
\langle AF \rangle_{bind} = \frac{\alpha'^2 |E_0^{inc}|^2 k^2}{4\pi \varepsilon d} \left[ k_z \cos(kd) \sin(k_z d) + k \cos(k_z d) \sin(kd) \right],
\]
\[
\langle BF \rangle_{bind} = -\langle AF \rangle_{bind}.
\]
In the case of \( \langle AF \rangle_{bind} \) the first term in the square brackets describes the interaction of the incident field at \( z_A \) with the dipole \( A \) induced by the wave scattered by the dipole \( B \) (induced by the incident wave at \( z_B \)). The second term in the bracket corresponds to the interaction between the dipole \( A \) placed at \( z_A \) (induced by the incident field) and the wave scattered by the dipole \( B \) placed at \( z_B \) (induced by the incident field therein). Both terms correspond to the second and third terms explained in (18b). Figure 2 illustrates this process in the case of single incident wave. This mechanism has been already described in more details for counter-propagating incoherent Bessel beams.\(^{10,20}\) As a result of this process the equilibrium positions of two optically bound nanoparticles are separated by a distance close to \( \lambda/2 \). However, the overall amplitude of this force decreases inversely with the inter-particle separation \( d \). Moreover this binding force depends on \( \alpha'^2 \) and it follows from (2) that the binding force raises with the sixth power of the particles radius.

By setting \( \phi = 0, \) the counter-propagating waves create a standing wave and the resulting forces (21) are
\[
\langle AF \rangle_{SWbind} = -\alpha' |E_0^{inc}|^2 \left\{ k_z \sin(2k_z z_A) - \frac{\alpha' k^2}{4\pi \varepsilon d} \left[ k_z \cos(kd) \sin(k_z d) + k \cos(k_z d) \sin(kd) \right] \right\},
\]
\[
\langle BF \rangle_{SWbind} = \alpha' |E_0^{inc}|^2 \left\{ -k_z \sin(2k_z z_B) - \frac{\alpha' k^2}{4\pi \varepsilon d} \left[ k_z \cos(kd) \sin(k_z d) + k \cos(k_z d) \sin(kd) \right] \right\}.
\]
Comparing to (22) it is seen that except the mutual distance between the particles \( d \), the force depends on the absolute position of one particle. These new terms (with respect to (22)) originate from the interference of
counter-propagating waves. The first term in curly brackets describes the force acting on a single nanoparticle placed in the standing wave. This problem has been already solved in more details for various types of standing waves - Gaussian, Bessel and evanescent. The first and second terms in the square brackets describe the optical binding described in (22). The third term in the square brackets describes the interaction between the incident field at \( z_A \) and the dipole \( A \) induced by the scattered radiation coming from the dipole \( B \) induced by the second counter-propagating incident wave. The fourth term corresponds to the interaction between the dipole \( A \) induced by the incident wave and the radiation scattered from the dipole \( B \) placed at \( z_B \) and induced by the counter-propagating wave.

The term in the front of the square bracket describes the influence of the mutual interaction between the particles with respect to the force coming from the incident standing wave. If we express this term using (2) we obtain

\[
\frac{m^2 - 1}{m^2 + 2} \frac{a^3 k^3}{kd^3}.
\]  

(25)

Therefore, if the size of the particle increases, the mutual interaction between the particles is stronger and can compete with the force coming from the incident standing wave. Unfortunately, the accepted approximation of induced dipoles ceased to be valid in this case.

2.4 Longitudinal binding of microparticles in counter-propagating interfering Bessel beams

We modified the coupled dipole method (CDM) to study numerically longitudinal optical binding of bigger particles. Here, we present new results related to the optical binding of bigger particles comparing to nanoparticles studied in the previous section. Figure 3 compares forces acting on the left particle (corresponds to the dipole \( A \) ) placed in incoherent or coherent Bessel beams (left or middle column) and on the right particle \( B \) placed in coherent Bessel beams. Five sizes of two identical particles are considered and shown in separate rows. In all cases it is assumed the left particle is fixed at \( z_A \) (equilibrium position in the incident standing wave) however the right particle moves along \( z \) axis. Figure proves that the mutual interaction between particles is stronger for bigger particles and also its amplitude is slightly stronger in coherent beams comparing to incoherent ones. For the biggest particle radius the presence of the second particle significantly shifts the oscillating force profile of the right particle.

3. EXPERIMENTAL OPTICAL BINDING OF MICROPARTICLES

In the experimental part we focused on the noticeable differences in the behaviour of particles in the interfering and non-interfering Bessel beams. We used the experimental setup shown in Figure 4. The fiber laser working at 1064 nm with maximal power 10 W was used. As the two beams passed through axicons they formed pseudo non-diffracting (Bessel) beams. Behind the axicon the wavevectors of refracted plane waves cover the surface of a cone with semi-apex angle \( \alpha_0 \) and the Bessel beam is a product of their interference. These beams went through a demagnifying telescope decreasing their high intensity Bessel beam core radius down to micrometer range. Bessel beam deformation caused by imperfections of the conical tips of the axicons were suppressed by centred beam-block placed in the Fourier plane (spatial filters SPF).

Colloidal sample consisted of polystyrene particles placed in an optical quality cuvette filled with water. The sample was diluted so that one particle crossed the beam within 2-5 minutes during quantitative measurements. Particle positions were observed in the scattered light and recorded by a fast CCD camera IDT XS-3. Thousands of frames were recorded with various concentrations of particles and number of bound objects in the beam. The individual particle positions were obtained from each frame by a correlation algorithm and interparticle separations were analyzed. We repeated the experiments in 1 mM NaCl solution with no observable change in the experimental results. Thus we ensured that any potential electrostatic interactions between particles have negligible effect at the observed interparticle separations. Figure 5 compares confinement of two polystyrene particles of radius 535 nm if both beams interfered (the first part of the record) or did not interfere (the second part of the record). When both beams interfered the particles were tightly localized in the standing wave structure of the incident beams and just three jumps occurred between neighbouring equilibrium positions of the bound structure. They were done by the neighbouring intensity maxima of the standing wave. Very well separated equilibrium positions are shown in the right plot of Figure 5.
Figure 3. System of two polystyrene particles of different radii (shown at different rows) placed on the optical axis at different interparticle distances $d$. Left particle is fixed with respect to the external optical field and the right particle moves along the optical axis. Optical force (blue line) and corresponding potential (green line) acting on nanoparticle in different configurations are computed using CDM: 1.col.: 2 counter-propagating incoherent Bessel beams & static left particle, 2.col.: 2 counter-propagating coherent Bessel beams & static and trapped (in SW optical pattern) left particle, 3.col.: 2 counter-propagating coherent Bessel beams & movable (in SW optical pattern) right particle. Polystyrene particles ($n_p = 1.59$) suspended in water ($n_m = 1.33$) and illuminated by a Bessel beam of vacuum wavelength 1064 nm, with $k_z/k = 0.99$ and the optical intensity at the centre of the Bessel beam equal to $0.63 \text{ mW/μm}^2$. 

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Figure 4. The experimental setup: LASER - low coherence fiber laser (IPG ILM-10-1064-LP), PBS - polarizing beam-splitter, metallic mirrors, axicons (cone lenses) with apex angle 170° (EKSPLA 130-0270), L1 - planoconvex lenses \( f = 150 \text{ mm} \) (Thorlabs LA1433-B), L2 - achromatic doublets \( f = 30 \text{ mm} \) for beam core radius 1.8 \( \mu \text{m} \), SPF - spatial filters, C - cuvette, LDO - Mitutoyo long working distance objective \( (f = 200 \text{ mm}, \text{numerical aperture 0.5, magnification 80}) \), CCD - fast CCD camera (IDT XS-3). The bottom insets show the cross-sections of optical intensity of both demagnified overlapping Bessel beams. The lengths of both paths were set and the halfwave plate was used to control the angle between both counter-propagating linearly polarized beams.

Figure 5. Left: Record of the measured separation between two polystyrene particles trapped in the overlapped cores of two counter-propagating Bessel beams. The first half of the record shows the case when both beams interfered and the particles were tightly localised in the incident standing wave. The second part shows the case when the polarization of one beam was rotated by \( \pi/2 \) and both beams did not interfere. Both particles kept bound but axial stiffness of their confinement is much lower. Right: Probability density of inter-particle separations calculated form the top record for both cases. Bessel beam core radius was \( \rho_0 = 1.8 \mu \text{m} \) and the incoming power was \( P = 6 \text{ W} \), the power in the Bessel beam core was 10 mW.
When the polarizations of the counter-propagating beams were mutually perpendicular, the particles were still optically bound but they were less tightly confined and consequently their interparticle distance changed more due to the thermal activation. The right plot of the probability density of interparticle distances proves this behaviour. Moreover it reveals several peaks separated by about half the laser wavelength in the medium. It is still open question if this modulation is a result of the binding between the particles or it is a residual interference due to the non-perfect polarization elements. Figure 6 shows similar results for the case of three bound particles.

Figure 6. The same record as in Figure 5 but for three optically bound particles. Presented data correspond to the distance between the left and middle particle.

The presented data correspond to the distance between the left and the middle particle. Again remarkable freezing of the structure in the interfering beams can be observed in the left figure.

4. CONCLUSIONS

We presented the theoretical model of the optical binding between two nanoparticles placed in two counter-propagating coherent linearly polarized beams. The polarization of one beam can be rotated from parallel to perpendicular orientation with respect to the second beam. This theory reveals that the optical binding becomes much stronger for bigger particles. We applied the coupled dipole method to express forces acting on two bigger particles and we verified the increase of the mutual interparticles interaction even for particles out of the applicability of the analytical model. Experimental results confirmed noticeable differences in the stiffnesses of the confined particles if they were bound in interfering beams or non-interfering beams. If the beams do not interfere, the particles keep to be optically bound but move axially over longer range due to the thermal activation. However, if the polarization is rotated to parallel, the positions of the particles are fixed with respect to the standing wave and their motion is much reduced. This simple experimental method can be used for controlled tuning of the inter-particle distances.

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