Laser Singular Theta-Pinch.

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The interaction of the two counter-propagating ultrashort laser pulses with a singular wavefronts in the thin slice of the underdense plasma is considered. It is shown that ion-acoustic wave is excited via Brillouin three-wave resonance by corkscrew interference pattern of a paraxial singular laser beams. The orbital angular momentum carried by light is transferred to plasma ion-acoustic vortex. The rotation of the density perturbations of electron fluid is the cause of helical current which produce the kilogauss axial quasi-static magnetic field. The exact analytical configurations are presented for an ion-acoustic current field and magnetic induction. The range of experimentally accessible parameters is evaluated.

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I. INTRODUCTION

The Stimulated Brillouin Scattering (SBS) in a plasma \[1,2\] is a subject of the considerable interest in a several recent decades \[3\]. The motivation of these research efforts is that a substantial power could be reflected from an underdense plasma via SBS thus limiting the absorption of laser radiation in an inertial confinement fusion (ICF) targets \[4\]. The possibilities of the ion-acoustic plasma wave resonant seeding by a controllable interference patterns of crossed laser beams \[5\] were studied in order to manipulate the SBS and Raman scattering \[6\]. On the other hand the SBS itself and in combination with a random phase plates (RPP) is considered as a possible tool for fusion beam smoothing and quenching of spatial instabilities in laser plasma \[7,8\]. The chaotic intensity and phase variations in a speckle pattern induced by RPP are also in use in wavefront reversal SBS mirrors in order to improve phase-conjugated (PC) replica fidelity \[9,10\]. As is shown recently the dark lines of a speckle (optical vortices) which are collocated with the optical phase singularities \[12\] inside SBS mirror volume are surrounded by the corkscrew interference patterns \[13,14\]. These optical helical patterns rotate synchronously with an acoustic angular frequency \(\Omega_\text{ia}\) and their rotation is accompanied by a circular motion of the SBS medium inside PC-mirror \[15\]. Such a circular motion of the medium carries an orbital angular momentum (OAM) \[15\] extracted from the exciting radiation \[16\]. In a similar way the transfer of OAM to an ensemble of \(10^4\) ultracold cesium atoms \[17\] via rotating interference pattern produced by a pair of the singular Laguerre-Gaussian(LG) laser beams had been observed in a nondegenerate four-wave mixing PC geometry. Quite recently the else example of an optical corkscrew patterns had been demonstrated with photorefractive Ba:TiO\(_3\) PC mirror \[18\].

The goal of this paper is in analysis of a new mechanism of an ion-acoustic wave (IAW) management by means of seeding the corkscrew interference pattern rotating around the isolated optical phase singularity. The qualitative picture is as follows. The optical helical pattern rotates due to a frequency detuning \(\Omega_p - \Omega_a\) between the pump and Stockes LG beams (fig.1) which collide inside freely expanding preformed underdense plasma jet (PJ, fig.2) created by evaporation of the solid-state target by nanosecond laser pulse. The parameters are expected to be close to experimental conditions summarized in \[7\]. The equilibrium atomic density \(n_{e0}\) is in the range \((0.2 - 0.07)\, n_c\) where \(n_c = 1.311 \cdot 10^{21}\, \text{cm}^{-3}\) is the critical density for laser wavelength \(\lambda = 0.5\, \mu\text{m}\) \[8\]. We suppose the detuning is adjusted in resonance with IAW having frequency \(\Omega_{ia} \approx 2\pi \cdot 10^{12}\) and fast damping rate \(\gamma_{ia} \approx (0.5 - 0.1)\, \Omega_{ia}\) \[8,10\]. In this SBS strong damping limit when duration of laser pulse is about \(\tau \approx 10^{-10}\, \text{sec}\) the helical density perturbations of the electron (\(\delta n_e\)) and ion (\(\delta n_i\)) liquids of the order of \(10^{-2}\, n_{e0}\) are collocated with the light intensity maxima. The rotation of the scalar density fields \(\delta n_{i,e}\) produces a nonrelativistic helical plasma flow with angular frequency \(\Omega_{ia}\). This helical current in the neutral plasma reminiscent to the geometry of a \(\Theta\) pinch \[1,2\] is expected to produce the axial quasi-static magnetic field. The mechanism discussed here is different from the process of the photon’s spin deposit from the circularly polarized intense light pulse to underdense plasma rotation via inverse Faraday effect described in \[20,21\] and already identified experimentally in \[22\] by measurement the megagauss (MG) axial magnetic induction \(B_z\). In our case the linearly polarized, i.e. zero spin LG photons \[23\] transfer their OAM to an ion-acoustic liquid via three-wave Brillouin resonance. The other difference is in the much lower (kilogauss range) magnitude of axial magnetic induction \(B_z\) resulting in the absence of the so-called IAW curtailment \[21\]. The resulting helical flow of electron liquid (fig.1) generates both axial and azimuthal magnetic inductions \(B_z\) and \(B_\theta\). We will show that interference pattern (fig.1) rotates as a ”solid body” and the vector of axial speed \(v_z\) is

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The most interesting issue is the interaction of electron and ion liquids in the process of SBS excitation by a short laser pulse having picosecond duration \((\tau \approx 10^{-10}-11\, \text{sec})\). The standard assumption of the plasma quasi-neutrality on the scales larger than Debye length \((r_D = \sqrt{k_B T_e e^2 n_e} \sim 10^{-6}-7\, \text{cm})\) requires the instantaneous compensation of the each electron density perturbation by an appropriate motion of the ion liquid. Thus within a framework of the conventional IAW model the macroscopic electron current must be exactly compensated by ionic current and net current ought to be equal to zero resulting in zero magnetic induction. This simple assumption is a basement of the three-wave SBS model (3 scalar ”parabolic” PDE) elaborated previously for the subnanosecond plasma flow with \(KeV\) temperatures [9]. Let us look more carefully to this assumption from the point of view of the optical pulse durations in the range mentioned above \((10 – 500\, \text{picoseconds})\): both the experiments followed by analytical estimates [10] and numerical simulations [9] supposed no magnetic fields. But the slight change of duration \(\tau\) down to the 1 picosecond revealed the substantial magnetic fields [24,25].

Despite the fact of a five order increase of the optical fluxes \((\text{up to } 10^{11}\, \text{W/cm}^2)\) requires the instantaneous appearance of a MG magnetic field for a 1\,\text{ps}\,\text{sec} flows requires proper interpretation. The possible interpretation is that a Debye screening do not cancel the net current completely due to partial imbalance in between the ion \((n_i)\) and electron \((n_e)\) liquids. As a result IAW is transformed into magnetosonic wave [1,20].

Noteworthy the other mechanisms of magnetic fields generation. For example the axial currents co-directed with picosecond laser pulse were identified experimentally by Faraday rotation polarimetry as a source of the toroidal magnetic field of the MG range [24,25]. Laser \(Z – \text{pinch}\) formed by return current of the hot electrons ejected from a thin copper wire by an ultrashort laser pulse was studied in [20]. The Weibel instability mechanism is responsible for generation of the quasistatic magnetic fields for the anisotropic electron’s density in both nonrelativitc helical plasma flows [27] and intense femtosecond laser pulse in an underdense plasma [28]. The singularities in density and magnetic field growing in time as \((t - t_0)^{-1/2}\) were considered earlier in [27]. The gigantic magnetic induction \(B\) mechanisms were discussed in [20,29,30].

II. HELICAL INTERFERENCE PATTERNS OF THE LAGUERRE-GAUSSIAN BEAMS AND ANGULAR MOMENTUM DENSITY

SBS is a decay of the quasimonochromatic electromagnetic wave \(E_p\) into the Stockes wave \(E_s\) and the ion-acoustic wave \(Q\) via Bragg scattering. The mechanism of this three-wave parametric instability [7] is the electrostrictive excitation of a sound wave by the moving interference pattern of two waves \((E_p\) and \(E_s\)) detuned by the sound frequency \(\Omega_{ia} = 2\Omega_p n_r c_{ia}/c [1,2,11]\), where \(c_{ia}\) is the velocity of the ion-acoustic wave, \(n_r\) is refractive index. The underdense plasma frequency \(\omega_p = \sqrt{n_e e^2/\epsilon_0 m_e}\) is assumed to be smaller than laser frequency \(\Omega_p \approx \omega_p/\sqrt{0.2}\).

In quantum terms the decay of the each pump photon to a Stockes photon and an ion-acoustic phonon follows to energy conservation condition \(h\Omega_p - h\Omega_s = h\Omega_{ia}\) and to momentum conservation \(\hbar \vec{k}_p - \hbar \vec{k}_s = \hbar \vec{k}_{ia} \approx 2\hbar k_p\) as well. The angular momentum conservation [13] also takes place resulting in OAM deposit in rotating electron liquid (rather than photon’s spin deposit [20]). In classical picture the detuning between optical waves \(\Omega_{ia}\) is a Doppler shift arising due to reflection from sound grating moving with the speed \(c_{ia}\) in a medium (plasma) with refractive index \(n_r\). For the backward SBS the period of sound grating \(p = \lambda_{ia}\) is approximately equal to a half of exciting wavelength \(\lambda_{ia} \approx \lambda_p/2\) (Bragg reflection condition). Noteworthy the Bragg and Doppler conditions are valid both near the bright spots and in the vicinities of the dark lines as well. The fundamental difference of the dark (vortex) lines from the bright spots of the speckle stems from the different structure of
their optical wavefronts. In a bright spot the wavefront is parabolic while near a vortex line the wavefront is helicoidal \[12, 13\]. In spite of the statistical nature of a speckle the complex amplitudes of the electric fields \(E_{p,s}\) for the counter-propagating waves could be written analytically inside a given bright or dark spot of a speckle. For a homogeneous polarization state the electric field has the following form in a cylindrical coordinate system \((z, r, \theta, t)\) nested in a given spot:

\[
E_{(p,s)}(z, r, \theta, t) = E_{0, p,s}(z, r) \cdot \exp \left\{ i \chi_{p,s}(z, r) \pm i \ell \theta \right\},
\]

(1)

where \(f_{p,s}(z, r)\) is a smooth amplitude function elongated in \(z\)-direction (e.g. Gaussian one), \(\chi_{p,s}(z, r)\) is a smooth phase profile (e.g. parabolic one), \(\ell\) is an azimuthal quantum number (topological charge). Define the mean square root transverse scale of a speckle entity as \(< D >\). Then mean longitudinal length of a variational speckle functions \(f_{p,s}(z, r)\) and \(\chi_{p,s}(z, r)\) is the Rayleigh range (Fresnel length) \(L_R \approx \sqrt{D^2/\lambda}\). Both \(f_{p,s}(z, r)\) and \(\chi_{p,s}(z, r)\) are assumed to be parabolic with respect to \(r\) near \(z\)-axis. For \(\ell = 0\) the equation (1) describes the phase-conjugation of a bright spot, e.g. zeroth-order Gaussian beam. For \(\ell \geq 1\) and when \(f_{p}(z, r) = f_{s}(z, r)\) and \(\chi_{p}(z, r) = \chi_{s}(z, r)\) the eq. (1) describes the perfect phase-conjugation of an optical vortex with ultimate fidelity i.e. when the correlation integral of the pump \(E_p\) and Stockes \(E_s\) fields is equal to unity \(\|E_p\|^2 = 1\). For the \(\ell\)-order \(LG_{\ell}^0\) beam (LG) the \(f_{p,s}(z, r)\) has Gaussian form \(\|E_p\|^2 = \exp \left( -r^2/\rho_c^2 \right)\). Thus two phase-conjugated LG compose in their common waist the corkscrew interference pattern (fig.1). Qualitatively the same corkscrew pattern in PC-mirror appears in the vicinity of each optical vortex line of an optical speckle pattern:

\[
|E_p(z, r, \theta, t) + E_s(z, r, \theta, t)|^2 = |E_p|^2 |1 + R_{pc} + 2 \sqrt{R_{pc}} \cos(\Omega p - \Omega s - k p z + 2 \ell \theta)| \times r^{2|\ell|} f_{pc}^2(z, r),
\]

(2)

where \(R_{pc} = |E_s|^2/|E_p|^2\) is the PC-reflectivity. The maximum of the interference pattern acts as an Archimedean screw transferring OAM from light to the matter \(\|E_p\|^2\). This happens because the orbital part of the electromagnetic angular momentum density \(M_z\) of LG is collocated with a maximum of the light intensity \(\|E_p\|^2\):

\[
M_z(z, r, \theta, t) \approx \frac{\ell}{\Omega p} |E_{p,s}|^2 + \frac{\sigma_z}{2 \Omega_p} \frac{\partial |E_{p,s}|^2}{\partial r},
\]

(3)

where \(\sigma_z = 0, \pm 1\) corresponds to linear, right and left polarizations of LG respectively. Hence there is a principal difference in between a vortex line and a bright spot in the optical speckle. The bright spot only suppresses the medium by rolls of interference pattern \(\|E_p\|^2\) along propagation direction \((z-axis)\) of a pump wave \(E_p\), while the vortex line located in a dark spot produces the additional rotational effect upon liquid and imprints the angular momentum therein \(\|E_p\|^2\).

III. OPTICAL COLLIDER SETUP FOR ULTRASHORT LASER PULSES WITH A SINGULAR WAVEFRONTS

The easiest way to produce the optical corkscrew by means of the interference of a pair of the two phase-conjugated optical vortices is to use a loop optical setup (fig.2) reminiscent to optical collider schemes \(\|E_p\|^2\). The Eq. (1) shows that a phase-conjugated optical vortex \(E_v \approx \exp(ik_v z - i \theta)\) has the same ratio of the signs of the \(k_\ell\) and \(\ell\) in a self-similar variable as those of a pump wave \(E_p \approx \exp(-ik_p z + i \theta)\). This means that a helical wavefront of the PC-vortex \(E_v\) has the same value and sign of the topological charge \(\ell\), as an incident wave \(E_p\). Hence the PC-vortex is absolutely identical to an incident wave except for the exactly opposite direction of propagation. The restriction on the turnover of the direction of propagation is a number of reflections from conventional, i.e. non-PC mirrors. The odd number of reflections changes the topological charge \(\ell\) to the opposite one \(\|E_p\|^2\). Hence the conceptual loop setup ought to perform the even number of reflections in order to keep the topological charges of colliding pulses identical to each other as shown at fig.2. The alternative is to use two mode converters (MC) at a final straight path before target PJ (also at fig.2). The other restriction imposed on coherence length means that a paths of the counter-propagating vortices ought to be equalized with an accuracy better than \(L_{coh} = c \tau\) for proper temporal overlapping.

For the transform-limited light pulses having \(\tau \approx 1 - 100 \text{ ps}\) duration the coherence length is \(L_{coh} \approx 3 \times 10^3 - 3 \times 10^3 \mu m\), i.e. above \(300 \cdot \lambda_p\). In such
a case the ultrashort pulses with a singular wavefronts divided by a beamsplitter in a proposed loop scheme (fig.2) will collide in a thin slice of a preformed underdense \((\Omega_{p} > \omega_{p} = \sqrt{n_{e}c^{2}/m_{e}\epsilon_{0}})\) plasma jet \([7]\) and produce corkscrew interference pattern therein. Because the SBS sound (IAW) grating parameters must be in a resonance defined by Doppler and Bragg conditions, the frequency shift between colliding pulses should be equal to the ion-acoustic frequency \(\Omega_{ia}\) with an accuracy better than ion-acoustic SBS linewidth \(\delta\Omega_{ia}\). Such frequency shift could be produced by means of Raman scattering in compressed gases.

The acoustical frequency \(\Omega_{ia}\) is of the order of \(10^{9}\) Hz for SBS in a room-temperature gases and liquids because it is determined by a thermal velocity \(V_{T} \approx \sqrt{k_{B}T/m_{e}}\) at an ambient temperature \(T \approx 300K (0.025 eV)\), where \(k_{B}\) is a Boltzmann constant and \(m_{e}\) is a mass of a particle. The Brillouin linewidth in transparent dielectric is defined by a damping rate of sound \([11]\). The ion-acoustic plasma frequency shift \(\Omega_{ia}\) has roughly the same dependence on electron and ion temperatures \(T_{e}\) and \(T_{i}\). In this case the much higher \(T_{i/c} \approx \, \gamma_{ia} \approx 1keV\) offers the \(10^{12}\) Hz frequency shift due to a basically same square root dependence \(c_{ia} \approx \sqrt{k_{B}T_{e}/m_{i}}\), or more precisely \(c_{ia} = \sqrt{k_{B}(Z_{e}/T_{e} + 3T_{i})/m_{i}}\), where \(Z\) is the ion’s charge \([9]\). The SBS linewidth for IAW is due to the Landau damping \(\gamma_{ia} \sim \nu_{L} = \Omega_{ia}/\sqrt{\pi Z m_{e}/8 m_{i}}\) \([33]\).

**IV. CONFIGURATION OF AN ION-ACOUSTIC PLASMA VORTEX SEED BY A CORSKREW INTERFERENCE PATTERN**

Below are the SBS equations of motion for the scalar slowly varying envelope optical fields, i.e. \(\mathbf{E}_{p}\) moving in the positive \(z\)-direction and \(\mathbf{E}_{s}\) moving oppositely:

\[
\frac{\partial \mathbf{E}_{p}(z, r, \phi, t)}{\partial z} + \frac{n_{e}}{c} \frac{\partial \mathbf{E}_{p}}{\partial t} + \frac{i}{2k_{p}} \nabla \cdot \mathbf{E}_{p} = i\frac{\Omega_{ia} n_{e}}{4 n_{e}c} \hat{Q} \mathbf{E}_{s} \tag{4}
\]

\[
\frac{\partial \mathbf{E}_{s}(z, r, \phi, t)}{\partial z} - \frac{n_{e}}{c} \frac{\partial \mathbf{E}_{s}}{\partial t} - \frac{i}{2k_{p}} \nabla \cdot \mathbf{E}_{s} = -\frac{\Omega_{ia} n_{e}}{4 n_{e}c} \mathbf{E}_{p} \hat{Q}^{*} \tag{5}
\]

and dimensionless slowly varying IAW density perturbation complex amplitude \(\hat{Q}\) \([12,13,33]\):

\[
\frac{\partial \hat{Q}(z, r, \phi, t)}{\partial z} + \frac{1}{\epsilon_{ia}} \frac{\partial \hat{Q}}{\partial t} + \frac{2\gamma_{ia} \hat{Q}}{\epsilon_{ia}} \frac{\partial}{\partial (k_{p} + k_{s})} \nabla \cdot \hat{Q} = \frac{i}{\epsilon_{0}} \frac{\epsilon_{ia}}{2 n_{e} k_{B} T_{e}} \tag{6}
\]

where \(\nabla = (\partial_{z}, \partial_{r}, \partial_{\phi})\), \(n_{e}\) is the critical plasma density \([4,9]\). For a preformed plasma jet of the thickness \(L_{jet} \approx 10^{4} \mu m\) \([5]\) probed by two colliding singular laser beams of comparable amplitudes \(\epsilon_{p} \approx \epsilon_{s}\) (fig.2) with beam waist diameter \(D \approx 100\mu m\) and wavelength \(\lambda_{p} \approx 1.05\mu m\) the Rayleigh range is \(L_{R} = k_{p}D^{2} \approx 10^{4}\mu m \approx 10L_{jet}\). Then \(\nabla \hat{Q}^{2}\) terms which are responsible for the tranverse effects could be omitted. In the linear regime of SBS and under the weak coupling conditions the fields \(\mathbf{E}_{p}\) and \(\mathbf{E}_{s}\) are close to those fixed at their free-space values at the opposite boundaries of the jet \(\mathbf{P}\) and their \(z\)-dependence is obtained exactly throughout all inner \(z\). This weak-coupling regime might be justified due to nonsaturated SBS reflectivity \((R_{sbs} \approx 0.63)\) obtained in \([5]\). Next assumption is the stationary SBS regime when IAW damping is assumed to be strong: \(\gamma_{ia} \approx (0.5-0.1)\Omega_{ia} > \tau^{-1} \approx 10^{10-11} Hz\) \([6,12]\). Thus exact expression for the complex amplitude of IAW \(Q(z, r, \theta, t) = \hat{Q} \exp[i(\Omega_{ia}t - k_{ia}z)]\) follows immediately for the pair of colliding LG vortices \([13]\) with \(|\epsilon_{p}^{0}| = |\epsilon_{s}^{0}|\):

\[
Q = \exp[i(\Omega_{ia}t - k_{ia}z) + i2\theta] \frac{r}{D} 2^{\mid t \mid} \\
\exp[- \frac{2r^{2}}{D^{2}(1 + z^{2}/(k_{p}D^{4}))}] \frac{\Omega_{ia}}{4n_{e}k_{B}T_{e}} |\epsilon_{p}^{0}|^{2}. \tag{7}
\]

Consequently the density perturbation field \(q(z, r, \theta, t) = \delta n/n_{e} = Re(Q)\) demonstrates a rotation typical to optical vortex, alike LG beam. The difference is in doubled topological charge \(2\ell\) due to the angular momentum conservation \([13,15]\):

\[
q = \cos[\Omega_{ia}t - k_{ia}z + 2\ell] \frac{r}{D} 2^{\mid t \mid} \\
\exp[- \frac{2r^{2}}{D^{2}(1 + z^{2}/(k_{p}D^{4}))}] \frac{\Omega_{ia}}{4n_{e}k_{B}T_{e}} |\epsilon_{p}^{0}|^{2}. \tag{8}
\]

Thus we see that in the colliding pulse geometry the IAW density perturbation \(q\) has a form of a double helix which rotates around \(z - axis\) with the angular frequency \(\Omega_{ia}\) (fig.1). Then an analytical expression for the ion-acoustic current density \(j_{ia}\) could be extracted from the dynamical equations for the electron liquid \([12,13]\):

\[
\frac{\partial q}{\partial t} + \nabla \cdot q \mathbf{V} = 0; \frac{d \mathbf{V}}{dt} = -\frac{e}{m_{e}} [\mathbf{E} + \mathbf{V} \times \mathbf{B}] \tag{9}
\]

From the first equation \([9](\text{continuity equation})\) the approximate expression for the ion-acoustic current density vector field \(j_{ia}(\mathbf{r}, t)\) may be obtained using following symmetry arguments. The \(q\) plays the role of the density of effective "charge" density of the long "charged twisted wire", while \(q \mathbf{V}\) is an effective "electric field" directed along the normal to this "charged twisted wire" (fig.1 \(x, y, z\) are unit vectors). Then the current density is obtained via Gauss theorem:

\[
\mathbf{j}_{ia}(\mathbf{r}, t) = 2\ell Z e \cdot n_{e} q(\mathbf{r}, t) \mathbf{V} = 2\ell Z e \cdot n_{e} q(\mathbf{r}, t) \times \mathbf{E}_{ia}(\mathbf{r}, t) \tag{10}
\]

Because the distribution functions \(W_{ne}(\mathbf{r}, \mathbf{p}, t)\) \([1]\) in the proposed geometry (fig.1) are expected to be strongly anisotropic and not known yet the usage of dynamical equations \([14,15]\) seems to be reasonable. Then substitution of the nonrelativistic \(j_{ia}\) in a Biot-Savarr integral gives the quasistationary magnetic field \(\mathbf{B}\):

\[
\mathbf{B}(\mathbf{R}, t) = \int \frac{\mu_{0} |j_{ia}(\mathbf{r}, t) \times (\mathbf{R} - \mathbf{r})|}{4\pi |\mathbf{R} - \mathbf{r}|^{3}} d^{3}\mathbf{r} \tag{11}
\]
The expression for magnetic induction inside the plasma solenoid (fig.1) and near it's end will be published elsewhere. In the far field the time-averaged induction \( \mathbf{B}(\mathbf{R}) = \tau^{-1} \int_0^\tau \mathbf{B}(\mathbf{R}, t) dt \) has the configuration of the magnetic dipole, whose sign is determined by the optical topological charge \( \ell \):

\[
\mathbf{B}(\mathbf{R}) = \frac{\mu_0}{4\pi} \left[ 3(\mathbf{\mu} \cdot \mathbf{R}) \mathbf{R} - \mathbf{\mu} \right]/R^5; \quad \mathbf{\mu} = 2\ell \frac{\delta n_{e,\parallel}}{\delta n_{e,\parallel}} \frac{I_{ia}}{\lambda_p} \frac{\pi D^2 L_{jet}}{\lambda_p},
\]

where \( \mathbf{\mu} \) is an effective magnetic moment induced by the corkscrew ion-acoustic current \( \mathbf{\mu} = 2\ell \frac{\delta n_{e,\parallel}}{\delta n_{e,\parallel}} \frac{I_{ia}}{\lambda_p} \frac{\pi D^2 L_{jet}}{\lambda_p} \).

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\]

This ratio proved to be of kinematic nature and this follows from a solid body rotation of the charged solenoid (fig.1): the slow axial current producing weak tangential field \( \mathbf{B}_{ia} \) is due to IAW translational motion while a fast tangential speed \( \Omega_{ia} \) produces the large axial induction \( \mathbf{B}_z \).

Apart from the isolated double helix interference pattern having diameter near \( D \approx 100\mu m \) in the proposed loop collider setup (fig.2), the more complex experimental situation appears in a speckle produced by RPP. The optical vortices in a RPP speckle have a size of \( \approx 10^4 \) to an azimuthal \( (|\mathbf{B}_\theta|) \) component of a static magnetic field is given by:

\[
\frac{|\mathbf{B}_z|}{|\mathbf{B}_\theta|} \approx \frac{e \cdot \delta n_e}{e \cdot \delta n_e} \frac{|v_{ia}|}{|v_z|} = \frac{\Omega_{ia} \cdot R}{v_{ia}} = \frac{2\pi D}{\lambda_p}.
\]

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In summary we have shown that the nonrelativistic plasma vortex flow resonantly initiated by SBS of the two phase-conjugated optical vortices is capable to produce a kilogauss quasistatic magnetic field in a thin slice of the preformed underdense plasma jet [7]. In accordance with our model (Eq. 11) the particles from a preformed thermal bath having KeV temperatures (e.g. \( T_i \approx 0.25 \) KeV, \( T_e \approx 0.4 \div 0.7 \) KeV [7]) are accelerated by rotating maxima of light intensity by virtue of ponderomotive force and acquire the azimuthal speed \( \Omega_{ia} \approx 0.1 < c \) in addition to axial speed \( v_{ia} \) which coincides with the IAW velocity. The azimuthal speed \( \Omega_{ia} \approx 0.1 < c \) is about two orders of magnitude compared to an axial one \( v_z \approx c_{ia} \) (13). This provides a corresponding increase of the axial component of magnetic induction \( \mathbf{B}_z \) compared to the azimuthal one \( \mathbf{B}_\theta \). The axial tesla-range static magnetic fields proposed here in rotating plasma microsolenoid setup might be interesting from the point of view of studies of the vacuum birefringence phenomena where a first results had been obtained quite recently in the PVLAS experiment [34] (noteworthy the different mutual orientation of the laser axis and magnetic induction in our case).

The case of relativistic intensities and femtosecond laser duration \( \tau \approx 10^{-14} \div 15 \) sec in a geometry (fig.2) was not studied yet and this case deserve a particular consideration. Nevertheless it is worth to mention the else interesting feature of the rotating interference pattern described by (2). For the sufficiently large radius \( R_{sl} \) the speed of the azimuthal motion \( \Omega_{ia} \) of the optical interference pattern may reach or even exceed the speed of light in a vacuum \( \Omega_{ia} R_{sl} \geq c \). It well known that such a superluminal motion of the interference maxima does not violate causality because formula (2) presumes the overlapping of infinitely long pulses. As is shown previously for the "fast light" phenomena [35] the interference maxima can not transfer the superluminal signal thus causality is maintained. The same situation holds for the nonlinear laser amplifiers [36, 37] and optical pulses in dispersive medium [38]. In our case when azimuthal speed of rotation \( \Omega_{ia} \) approaches \( c \), the actual relativistic plasma flow will be different from the above formulated nonrelativistic SBS model which presumes the collocation of the optical intensity maxima and plasma current.

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[1] A.F.Alexandrov, L.S.Bogdankevich and A.A.Rukhadze, Principles of plasma electrodynamics, Springer-Verlag,
Berlin (1984).

[2] W.L. Kruer , *The Physics of Laser Plasma Interactions*, Addison-Wesley, New York (1990).

[3] C.S.Liu, M.N.Rosenbluth, R.B.White, Phys.Fluids 17, 1211 (1974).

[4] P.M.Lushnikov and H.A.Rose, Phys.Rev.Lett. 92, 255003 (2004).

[5] H.A.Baldis, C.Labaune, E.Schifano, N.Renard, and A.Michard, Phys.Rev.Lett. 77, 2597 (1996).

[6] R.K.Kirkwood et al., Phys.Plasmas, 4, 1800 (1997).

[7] C.Labaune1, H.A.Baldis, and V.T.Tikhonchuk, Europhys.Lett., 38, 31 (1997).

[8] A.V.Maximov, I.G.Ourdev, D.Pesme, W.Rozmus, V.T.Tikhonchuk and C.E.Capjack, Phys.Plasmas, 8, 1319 (2001).

[9] P.Loiseau et al., Phys.Rev.Lett. 97, 205001 (2006).

[10] N.G.Basov, I.G.Zubarev, A.B.Mironov, S.I.Mikhailov and A.Y.OKulov, JETP, 52, 847 (1980).

[11] B.Y.Zeldovich, N.F.Pilipetsky and V.V.Shkunov, *Principles of Phase Conjugation*, (Berlin:Springer-Verlag)(1985).

[12] J.F.Nye and M.V.Berry, Proc. R. Soc. London, Ser., A336, p.165 (1974).

[13] A.Yu.Okulov, J.Phys.B., 41, 101001 (2008).

[14] A.Yu.Okulov, Phys.Rev.A., 80, 163907 (2009).

[15] A.Yu.Okulov, Phys.Rev.Lett., 88, 631 (2008).

[16] R.Marchiano, F.Coulouvrat, L.Ganjehi and J.L.Thomas, Phys. Rev. E 77, 016605 (2008).

[17] J.W.R.Tabosa and D.V.Petrov, Phys. Rev. Lett. 83, 4967 (1999).

[18] M.Woerdemann, C. Alpmann and C.Denz Optics Express, 17, 22791 (2009).

[19] V.N.Tsytovich, Phys. Usp. 177, 428 (2007).

[20] M.G. Haines Phys.Rev.Lett. 87, 135005 (2001).

[21] V.Yu. Bychenkov , V.I. Demin and V.T.Tikhonchuk, JETP, 78, 62 (1994).

[22] Z. Najmudin , et al. Phys.Rev.Lett. 87, 215004 (2001).

[23] J.Leach , M.J.Padgett, S.M.Barnett, S.Franke-Arnold, J.Courtial, Phys.Rev.Lett. 88, 257901 (2002).

[24] M.Borghesi, A.J.Mackinnon, A.R.Bell, R.Gaillard, O.Will, Phys.Rev.Lett. 80, 112 (1998).

[25] M.Borghesi, A. J. Mackinnon, R. Gaillard, O. Will, A. Pukhov, and J. Meyer-ter-Vehn, Phys.Rev.Lett. 81, 5137(1998).

[26] F.N.Beg et al., Phys.Rev.Lett. 92, 095001 (2004).

[27] V.Yu. Bychenkov, V.P.Silin and V.T.Tikhonchuk, JETP, 71, 709 (1990).

[28] V.P.Krainov, J.Phys.B., 36, 3187 (2003), V.P.Krainov, JETP, 96, 430 (2003).

[29] R.N.Sudan, Phys.Rev.Lett. 70, 3075 (1993).

[30] V.S.Belyaev, V.P.Krainov, V.S.Lisitsa, A.P.Matafonov, Phys.Usp. 51, 428 (2008).

[31] L.Allen, M.W.Beijersbergen, R.J.C.Spreeuw and J.P.Woerdman, Phys.Rev.A, 45, 8185 (1992).

[32] G.A.Mourou, T.Tajima and S.V.Bulanov, Rev.Mod.Phys., 78, 309 (2006).

[33] E.M.Lifshitz and L.P.Pitaevskii, *Physical Kinetics* (Landau and Lifshitz Course of Theoretical Physics) Butterworth-Heinemann, Oxford (1982).

[34] M. Marklund and P.K.Shukla , Rev.Mod.Phys., 78, 591640 (2006).

[35] B.M.Bokotovskii and A.V.Serov, Phys.Usp., 48(9), 903 (2005).

[36] R.V.Ambartsumyan, N.G.Basov, V.S.Letokhov et al., JETP, 23, 16 (1966).

[37] A.Yu.Okulov and A.N.Oraevskii, Sov.J.Quantum Electron. 18(2), 233 (1988).

[38] R.Y.Chiao, Phys.Rev.A., 50, R34 (1993).

[39] A.Yu.Okulov, private communication, (2010).