

# Enhanced direct laser vacuum acceleration of a charged particle in crossing plane-wave laser beams

H. Lin, C.P. Liu, C. Wang, B.F. Shen

**Abstract.** Strict theory shows that in the field of two laser beams crossing in vacuum with a phase delay  $\pi/2$ , a unique dynamics of the electron is observed. Unlike a single-beam case in which the velocity component along the laser magnetic field is always zero, the electron motion in the case of crossing beams is completely three-dimensional and hence a larger maximum value of the relativistic factor  $\Gamma$ , i.e. acceleration of an electron to larger energies, can be achieved.

**Keywords:** laser acceleration of a charged particle, laser beams with plane wavefronts.

## 1. Introduction

Direct vacuum acceleration of a charged particle by the laser field attracted considerable attention about ten years ago. Even though experimental [1–5] and theoretical [6–23] investigations outlined the physics of this issue, there are still strong doubts about the correctness of the results obtained [24, 25]. The criticism is mainly based on some theoretical assumptions of previously developed theories. For example, the exact form of the driving laser is questioned by McDonald [24]. Moreover, some authors focus their attention on the inhomogeneous laser beam shape and hence stress the role of the axial laser electric field (which appears in vacuum as a result of the condition  $\nabla \cdot \mathbf{E} = 0$ ) during acceleration [7, 14]. However, it is still difficult to understand how such an axial component of the laser vector potential appears in vacuum: If before entering the vacuum a laser beam does not have such an axial component when it travels in a medium, its descendant, or a laser beam in vacuum, will also have no axial component [26].

We temporarily ignore different viewpoints on the axial laser electric field in vacuum. The first question to be answered is whether or not laser radiation with a homogeneous transverse cross section of the field distribution can accelerate electrons. Note that such laser radiation in vacuum can have no axial laser electric field. Actually, the authors of [12] found out that an electron in a laser beam has an oscillating axial velocity. Maybe the authors were more interested in a shaped (or focused) laser beam, and therefore more in-depth investigations on the case of a shapeless laser beam (or a laser pulse

which is not focused and hence has a transverse homogeneous intensity profile) were not included in their works [7–15, 18].

At present, the main attention of researchers is attracted to laser-plasma based acceleration due to setting up a high-voltage electrostatic structure within plasmas by using a high-intensity laser beam [27]. This high-voltage electrostatic structure acts as conventional solid-state accelerator but works at higher acceleration gradients. In this case, the amplitude of such a high-voltage electrostatic structure is proportional to laser intensity, which leads to a necessity to enhance the laser intensity [28].

Recent investigations have shown that a laser beam without transverse intensity modulation can cause oscillation of the axial velocity of an electron with an amplitude whose value is insensitive to laser intensity [29]. This confirms the feasibility of vacuum acceleration by such a shapeless laser beam. According to Ref. [29], the axial velocity of the electron oscillates periodically at a frequency comparable to the laser one. Such a periodic oscillation is due to a driving laser beam, which causes the electron motion being described mathematically by the first integral. This enlightens us to consider what will happen if we apply a more complicated configuration of the driving beams. Esarey et al. [14] have considered the configuration of two crossing shaped beams and found that the acceleration in this configuration depends on the axial component of each laser beam. Here, we study the configuration of two crossing shapeless beams in the absence of the axial electric field. Some authors also note the considerable effect of the configuration of two interfering fields on the electron acceleration [30, 31].

Usually the electron transverse velocity  $v_{\perp}$  is assumed equal to  $eA/m_e$  ( $A$  is the vector potential of the laser radiation field) [12]. However, if  $v$  is initially zero,  $A$  is also zero. Thus,  $A$  depends on time as function  $\sin$  and hence laser electric field  $E$  – as  $\cos$ , which is not initially equal to zero. Because the laser energy density is proportional to  $E^2$ , the boundary condition  $E|_{t=0} = 0$  is more realistic to describe the laser field, and therefore,  $E$  should depend on time as  $\sin$  and hence  $A$  – as  $\cos$ . This means that  $v_{\perp} = eA/m_e - eA|_{t=0}/m_e$ . In this case  $v_{\perp}$  periodically changes its direction and hence the electron is bound to the laser beam axis, unlike the case  $v_{\perp} = eA/m_e$  when  $v_{\perp}$  periodically changes its value and the electron can leave the laser beam axis. If the laser beam has a finite transverse size, the electron can entirely escape the laser beam.

Such a ‘side’ escape is of practical value. Some authors have proposed a vacuum acceleration mechanism based on a laser pulse of noninteger cycle [32–35], which gives longitudinal acceleration to electrons; however, this pulse is difficult to generate. Nevertheless, even for a laser pulse of integer

H. Lin, C.P. Liu, C. Wang, B.F. Shen State Key Laboratory of High Field Laser Physics, Shanghai Institute of Optics and Fine Mechanics, P. O. Box 800-211, Shanghai 201800, China; e-mail: linhai@siom.ac.cn

Received 17 April 2014; revision received 12 October 2014  
Kvantovaya Elektronika 45 (9) 863–867 (2015)  
Submitted in English

cycle, from which electrons can escape, there appears a longitudinal force felt by the escaped electrons, and the duration of its effect of such electrons can be equal to an odd number of periods.

Therefore, the authors of Ref. [29] have concluded that a laser beam of finite transverse size can transport and accelerate free electrons. However, such acceleration is limited because the electron velocity periodically oscillates in some range. Here, we study whether or not multiple laser beams can lead to better vacuum acceleration of charged particles.

## 2. Theory

First we write electric/magnetic field components in three directions

$$\begin{array}{ccc} e_x & e_y & e_z \\ E_0 \sin(kz - \omega t) & E_0 \cos(kx - \omega t) & 0, \\ 0 & B_0 \sin(kz - \omega t) & B_0 \cos(kx - \omega t). \end{array}$$

In such a configuration, two driving laser beams with a phase delay  $\pi/2$  are directed along the  $z$  and  $x$  axes, respectively. The dynamics of a classic single-body charge in a monochromatic laser beam has been rigorously described by some authors [6, 18, 36, 37]. A pronounced feature of the single-beam case is that the motion along the laser magnetic field is forbidden. Apparently, the motion in the crossing-beam case is completely three-dimensional (3D).

We start with relativistic Newton equations:

$$F_x = E_0 \sin(kz - \omega t) + \left[ -\frac{dz}{dt} B_0 \sin(kz - \omega t) + \frac{dy}{dt} B_0 \cos(kx - \omega t) \right] = \Gamma \frac{d^2 x}{dt^2} + \frac{dx}{dt} \frac{d\Gamma}{dt}, \quad (1)$$

$$F_y = E_0 \cos(kx - \omega t) + \left[ 0 - \frac{dx}{dt} B_0 \cos(kx - \omega t) \right] = \Gamma \frac{d^2 y}{dt^2} + \frac{dy}{dt} \frac{d\Gamma}{dt}, \quad (2)$$

$$F_z = 0 + \left[ \frac{dx}{dt} B_0 \sin(kz - \omega t) - 0 \right] = \Gamma \frac{d^2 z}{dt^2} + \frac{dz}{dt} \frac{d\Gamma}{dt}. \quad (3)$$

Here,  $F$  is the Lorentz force; the  $e/m_e$  is included in  $E_0$  and  $B_0$ ; and  $\Gamma$  is the relativistic factor of the electron. Because for light in vacuum  $E_0 = cB_0$ , we can rewrite Eqns (1–3) as

$$\frac{d}{ds} \left[ \Gamma \left( \frac{d\varphi}{ds} + 1 \right) \right] = \frac{1}{c} \frac{dy}{dt} \frac{\omega_B}{\omega} \cos \varphi - \frac{\omega_B}{\omega} \sin \theta \frac{d\theta}{ds}, \quad (4)$$

$$\frac{d}{dt} \left( \Gamma \frac{1}{c} \frac{dy}{dt} \right) = -\frac{\omega_B}{\omega} \frac{d \sin \varphi}{dt}, \quad (5)$$

$$\frac{d}{ds} \left[ \Gamma \left( \frac{d\theta}{ds} + 1 \right) \right] = \left( \frac{d\varphi}{ds} + 1 \right) \frac{\omega_B}{\omega} \sin \theta, \quad (6)$$

where  $s = kz - \omega t$ ,  $\varphi = kx - \omega t$ ,  $kc = \omega$ ,  $s = \omega t$ ,  $(dz/dt)c^{-1} = d\theta/ds + 1$ ,  $(dx/dt)c^{-1} = d\varphi/ds + 1$  and  $\omega_B = B_0$ . It follows from Eqn (5) that

$$\frac{1}{c} \frac{dy}{dt} = -\frac{\omega_B}{\omega} \frac{\sin \varphi}{\Gamma}, \quad (7)$$

which leads to a new expression for  $\Gamma$

$$\Gamma^2 = \frac{1 + (\omega_B/\omega)^2 \sin^2 \varphi}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2}. \quad (8)$$

Thus, Eqns (1–3) are reduced into coupled second-order nonlinear differential equations of  $\theta$  and  $\varphi$

$$\frac{d}{ds} \left[ \Gamma \left( \frac{d\varphi}{ds} + 1 \right) \right] = -\left( \frac{\omega_B}{\omega} \right)^2 \sin \varphi \cos \varphi - \frac{\omega_B}{\omega} \sin \theta \frac{d\theta}{ds}, \quad (9)$$

$$\frac{d}{ds} \left[ \Gamma \left( \frac{d\theta}{ds} + 1 \right) \right] = \left( \frac{d\varphi}{ds} + 1 \right) \frac{\omega_B}{\omega} \sin \theta. \quad (10)$$

Note that Eqns (9), (10) belong to following general form of differential equations

$$C_1 \frac{d^2 \varphi}{ds^2} + C_2 \frac{d^2 \theta}{ds^2} = D_{12}, \quad (11)$$

$$C_3 \frac{d^2 \varphi}{ds^2} + C_4 \frac{d^2 \theta}{ds^2} = D_{34}. \quad (12)$$

Here,  $C_{1,2,3,4}$  and  $D_{1,2,3,4}$  are the functions of  $(\theta, \varphi, d\theta/ds, d\varphi/ds)$ :

$$\begin{aligned} C_1 &= \left[ \frac{1 + (\omega_B/\omega)^2 \sin^2 \varphi}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \right]^{1/2} \\ &\quad \times \left[ \frac{1 - (d\theta/ds + 1)^2}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \right], \\ C_2 &= \left[ \frac{1 + (\omega_B/\omega)^2 \sin^2 \varphi}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \right]^{1/2} \\ &\quad \times \left[ \frac{(d\theta/ds + 1)(d\varphi/ds + 1)}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \right], \\ C_3 &= \left[ \frac{1 + (\omega_B/\omega)^2 \sin^2 \varphi}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \right]^{1/2} \\ &\quad \times \left[ \frac{(d\theta/ds + 1)(d\varphi/ds + 1)}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \right], \end{aligned} \quad (13)$$

$$\begin{aligned} C_4 &= \left[ \frac{1 + (\omega_B/\omega)^2 \sin^2 \varphi}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \right]^{1/2} \\ &\quad \times \left[ \frac{1 - (d\varphi/ds + 1)^2}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \right], \end{aligned}$$

$$\begin{aligned} D_{12} &= -\left( \frac{\omega_B}{\omega} \right)^2 \sin \varphi \cos \varphi - \frac{\omega_B}{\omega} \sin \theta \frac{d\theta}{ds} \\ &\quad - \frac{1}{2} \left[ \frac{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2}{1 + (\omega_B/\omega)^2 \sin^2 \varphi} \right]^{1/2} \\ &\quad \times \frac{(\omega_B/\omega)^2 \sin(2\varphi)(d\varphi/ds)}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \frac{d\varphi}{ds}, \end{aligned}$$

$$D_{34} = \left( \frac{d\varphi}{ds} + 1 \right) \frac{\omega_B}{\omega} \sin \theta$$

$$- \frac{1}{2} \left[ \frac{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2}{1 + (\omega_B/\omega)^2 \sin^2 \varphi} \right]^{1/2}$$

$$\times \frac{(\omega_B/\omega)^2 \sin(2\varphi) (d\varphi/ds)}{1 - (d\theta/ds + 1)^2 - (d\varphi/ds + 1)^2} \frac{d\theta}{ds}.$$

Mathematically, because  $\Gamma$  is determined by Eqns (9), (10) (or by two variables  $\theta$  and  $\varphi$ ), there are two variants of the evolution of  $\Gamma$ . At a continuous growth of  $\Gamma$  [ $d\Gamma/dt = \sum E_{x,y,z}(dx,y,z/dt)$ ], its maximum increment can be estimated as  $\Delta\Gamma = \int E(r(t), t) v(t) dt \leq \int |E(r(t), t)| |v(t)| dt < \int |E(r(t), t)| c dt \leq \int |E|_{\max} c dt = |E|_{\max} ct = \omega A_{\max} ct$ . It follows from this estimate that for large  $\Gamma$  to be obtained, a large peak intensity of laser radiation is needed.

At some combinations of the parameter  $\theta$ ,  $\varphi$ ,  $d\theta/ds$ ,  $d\varphi/ds$ , defined by

$$0 = 1 - \left( \frac{d\theta}{ds} + 1 \right)^2 - \left( \frac{d\varphi}{ds} + 1 \right)^2, \quad (14)$$

the condition  $C_1 C_4 - C_2 C_3 = 0$  exists and, therefore,  $|d^2\varphi/ds^2| = |d^2\theta/ds^2| = \infty$  if only the condition  $D_{12} = 0 = D_{34}$  is not met. The existences of these special combinations suggest a singular (or discontinuous) growth of  $\Gamma$  at which there is sudden jump till the value, significantly exceeding the maximum ones at a continuous growth of this parameter.

Except these points, only the continuous evolution needs to be taken into account. An overall estimation by  $\Gamma$  should contain both ways. Indeed, if  $C_1 C_4 - C_2 C_3 = 0$ , then  $\lim_{E \rightarrow 0} (d\Gamma/dt) \rightarrow 0$ , while in the opposite case,  $\lim_{E \rightarrow 0} (d\Gamma/dt) \rightarrow 0$ . However, in accordance with (14), the first variant of the evolution takes place at  $\Gamma = \infty$ .

We should also note that the electric field is absent along the  $z$  direction ( $E_z = 0$ ). Thus, the formula  $d\Gamma/dt = \sum E_{x,y}(dx,y/dt)$  only reflects the time-variation of  $\Gamma$  when  $d^2z/dt^2 = 0$  (i.e.,  $dz/dt$  is independent of time), and  $d\Gamma/dt | dz/dt = \text{const} = \sum E_{x,y}(dx,y/dt)$ . If  $dz/dt = \text{const}$ , Eqn (3) can be reduced to the expression  $d(\Gamma dz/dt)/dt = (dx/dt) B_0 \sin(kz - \omega t)$ , which does not give a contribution to  $d\Gamma/dt$ .

The quantity  $d\Gamma/ds$  consists of three parts: one is proportional to  $d^2\varphi/ds^2$ , the second – to  $d^2\theta/ds^2$  and the third is due to  $d[1 + (\omega_B/\omega)^2 \sin^2 \varphi]/ds$  and independent of either  $d^2\varphi/ds^2$  or  $d^2\theta/ds^2$ . The  $d^2\theta/ds^2$ -independent part can be described by the formula  $d\Gamma/dt | d^2\theta/ds^2 = 0 = \sum E_{x,y}(dx,y/dt)$ . Clearly, it is insufficient to calculate  $\Gamma$  merely from the formula  $d\Gamma/dt = \sum E_{x,y}(dx,y/dt)$ . To calculate the evolution of  $\Gamma$ , we should calculate  $d^2\theta/ds^2$  and  $d^2\varphi/ds^2$  from Eqns (11), (12) and then calculate  $\Gamma$  at a next time point. The presence of the condition  $d^2\theta/ds^2 \neq 0$  is the reason of the dependence  $\Gamma(t)$ , differing from  $\Gamma(t) = \Gamma(t=0) + \int_0^t [\sum E_{x,y}(dx,y/dt)] dt = 0$ .

### 3. Numerical results and discussions

For the same set of parameters ( $d\theta/ds$ ,  $d\varphi/ds$ ), according to Eqn (8),  $\Gamma$  in the 3D case is formally larger than that in the 2D case ( $\Gamma_{2D} = \sqrt{1 + (\omega_B/\omega)^2 \sin^2 \varphi}^{-1}$ ). This implies that under same field strength of the driving laser, a higher maximum value of  $\Gamma$  can be achieved in the 3D case. Note that if

each beam has a peak field strength  $(\sqrt{2}/2) eA_{\max} (m_e c^2)^{-1}$ , the energy contained in two beams is equal to that in a single beam whose peak strength is  $eA_{\max} (m_e c^2)^{-1}$ . Thus, as shown in Fig.1, the energy contained in two beams at  $\omega_B/\omega = 1.5$  is equal to that in a single beam at  $\omega_B/\omega = 1.5\sqrt{2} = 2.2$ . Clearly, according to a familiar estimation, the maximum value of  $\Gamma$  at  $\omega_B/\omega = 2.2$  is  $\Gamma_{\max} = \sqrt{1 + 2.2^2} = 2.37$ , and in two beams at  $\omega_B/\omega = 1.5$  it can reach 10. At  $\omega_B/\omega = 2.5$ ,  $\Gamma_{\max}$  in two beams can be up to 70 or higher, while in one beam at  $\omega_B/\omega = 2.5\sqrt{2} = 3.6$ ,  $\Gamma_{\max} = 4$ .

Because a short-wavelength light source (X-ray or shorter wavelength) is significant for many applications, we are also interested in whether or not the studied configuration can act as an effective source of short-wavelength dipole radiations. Data presented in Fig. 1c suggest that the studied configuration is not an ideal short-wavelength light source because of predominant low-order harmonics.

As previously pointed out, in the single-beam case either the condition  $v_{\perp} = eA/m_e - eA|_{t=0}/m_e$  or  $p_{\perp} = eA - eA|_{t=0}$  is fulfilled. This relation, or its approximated counterpart  $p_{\perp} = eA$  allows one to calculate exactly  $\Gamma$ . However, in the case of multiple beams from different directions, this relation is violated. Indeed, the presence of the term  $(dy/dt) B_0 \cos(kx - \omega t)$  in (1) implies that  $d^2x/dt^2 = dA_x/dt = E_0 \sin(kz - \omega t) + [-(dz/dt) \times B_0 \sin(kz - \omega t)]$ . This means that the condition  $v_{\perp} = eA/m_e - eA|_{t=0}/m_e$  or  $v_{\perp} = eA/m_e$  does not hold. Nevertheless the known relation  $\Gamma = \sqrt{1 + (e/m_e c)^2 A^2}$  is valid for the single-beam case, which indicates that a higher  $\Gamma$  can be obtained by increasing  $A^2$ .

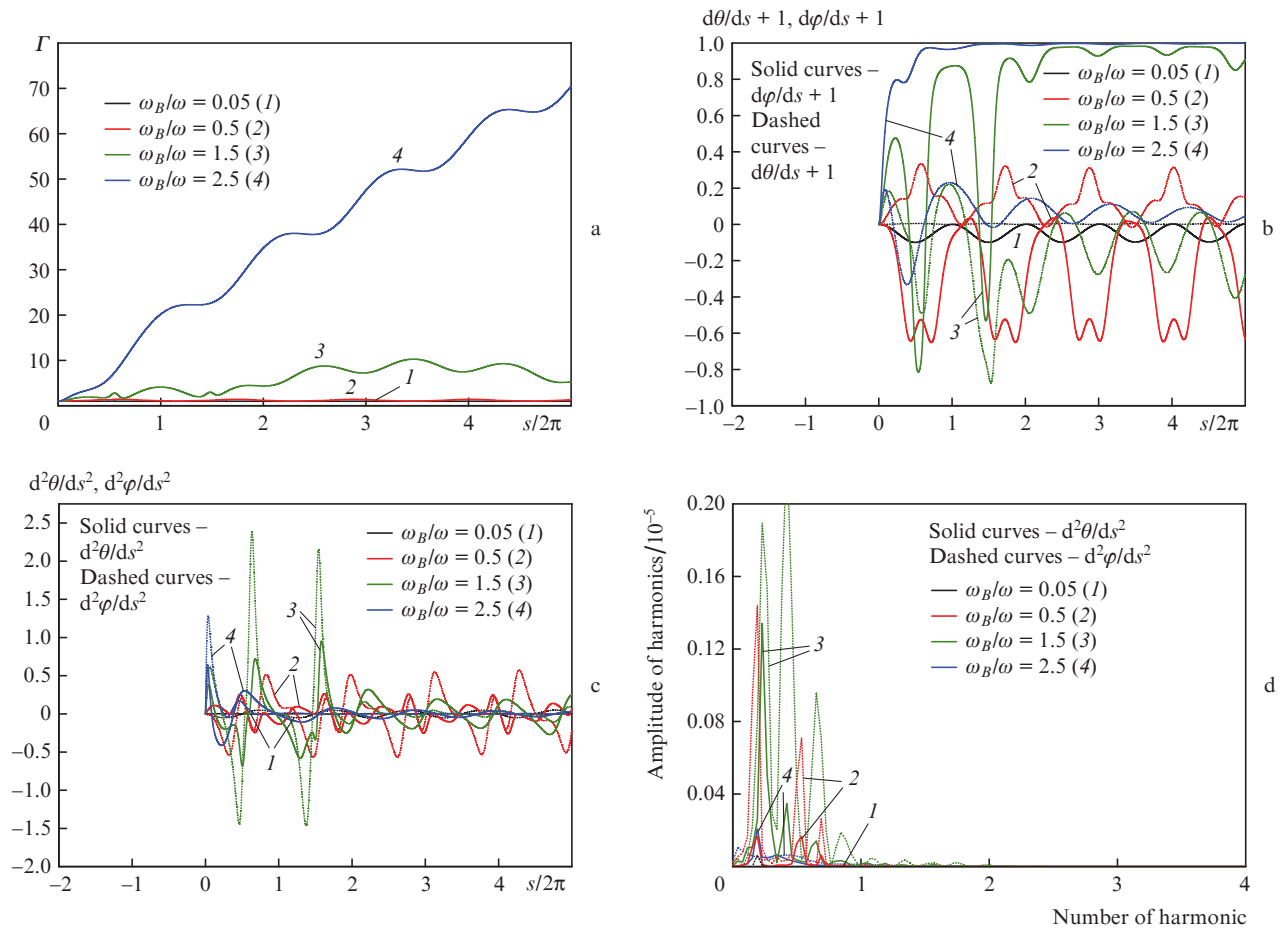
Indeed, in the single-beam case the field is described by the following components:

$$\begin{array}{ccc} e_x & e_y & e_z, \\ E_0 \sin(kz - \omega t) & 0 & 0, \\ 0 & B_0 \sin(kz - \omega t) & 0. \end{array}$$

In this case, the relations  $(dy/dt)c^{-1} = 0$  and  $(dy/dt)c^{-1} = 1 - \Gamma^{-1}$  are valid, which yield  $[(dx/dt)c^{-1}]^2 = 2[1 - (dz/dt)c^{-1}](dz/dt)c^{-1}$ . Therefore, the dynamics in the single-beam case is completely 1D (i.e., is determined by only one variable  $\theta$ ) and described by the expression  $(dz/dt)c^{-1} = d\theta/ds + 1$ . In addition, the classic orbit is open in both polarisation directions and the propagation of radiation is possible only in one direction (because the inequality  $(dz/dt)c^{-1} \geq 0$  must hold in order to ensure the fulfilment of the condition  $2[1 - (dz/dt)c^{-1}] \times (dz/dt)c^{-1} \geq 0$  [36, 37]). In the single-beam case, there is only the second-order nonlinear differential equations of  $\theta$ :  $C_4 d^2\theta/ds^2 = D_4$  [where detailed expressions for  $C_4$  and  $D_4$  can differ from those in Eqn (13) but do not affect the validity of following conclusion], but the equality  $C_4 = 0$  will only occurs when  $1 = d\theta/ds + 1 = (dz/dt)c^{-1}$ . This causes a discontinuous growth of  $\Gamma$ , which is initially absent in the single-beam case.

### 4. Conclusions

Strict theory and numerical results show that a complete 3D motion of a classic charge can be achieved in the case of two crossing shapeless laser beams with a phase delay  $\pi/2$ . Unlike the single-beam case in which the motion is two-dimensional, the maximum value of  $\Gamma$  for completely 3D motion can be higher under same laser field strength. In addition, in the case of crossing-beams with a phase delay  $\pi/2$ , larger temporal oscillations of charge acceleration lead to the generation of



**Figure 1.** Evolution of (a)  $\Gamma$ , (b) velocity and (c) acceleration rate, as well as (d) its Fourier spectrum.

radiation at high-order harmonics, which can be used for creating a high-frequency photon source.

**Acknowledgements.** This work is supported by the National Natural Science Foundation of China (Grant Nos 11374318, 11335013 and 11174305).

## References

- Moore C.I., Knauer J.P., Meyerhofer D.D. *Phys. Rev. Lett.*, **74**, 2439 (1995).
- Malka G., Lefebvre E., Miquel J.L. *Phys. Rev. Lett.*, **78**, 3314 (1997).
- Kawata S., Maruyama T., Watanabe H., Takahashi I. *Phys. Rev. Lett.*, **66**, 2072 (1991).
- Hussein M.S., Pato M.P. *Phys. Rev. Lett.*, **68**, 1136 (1992).
- Hussein M.S., Pato M.P., Kerman A.K. *Phys. Rev. A*, **46**, 3562 (1992).
- Landau L.D., Lifshitz E.M. *The Classical Theory of Fields* (London: Pergamon Press, 1971).
- Lax M., Louisell W., McKnight W. *Phys. Rev. A*, **11**, 1365 (1975).
- Steinhauer L.C., Kimura W.D. *J. Appl. Phys.*, **72**, 3237 (1992).
- Hauser T., Scheid W., Hora H. In: *Laser Interaction and Related Plasma Phenomena*. Ed. by G.H. Miley (New York: AIP, 1994) p. 20.
- Haaland C.M. *Opt. Commun.*, **114**, 280 (1995).
- Sprangle P., Esarey E., Krall J., Ting A. *Opt. Commun.* (in press).
- Hartemann F.V. et al. *Phys. Rev. E*, **51**, 4833 (1995); *ibid.* **58**, 5001 (1998).
- Troha A.L. et al. *Phys. Rev. E*, **60**, 926 (1999).
- Esarey E., Sprangle P., Krall J. *Phys. Rev. E*, **52**, 5443 (1995).
- Esarey E., Sprangle P., Pillot F.M., Krall J. *J. Opt. Soc. Am. B*, **12**, 1695 (1995).
- Brown L.S., Kibble T.W.B. *Phys. Rev.*, **133**, A705 (1964).
- Sarachik E.S., Schappert G.T. *Phys. Rev. D*, **1**, 2738 (1970).
- Bochovne E.J., Moore G.T., Scully M.O. *Phys. Rev. A*, **46**, 6640 (1992); Scully M.O., Zubairy M.S. *ibid.* **44**, 2656 (1991).
- Corkum P.B., Burnett N.H., Brunel F. *Phys. Rev. Lett.*, **62**, 1259 (1989).
- Andreev S.N., Makarov V.P., Rukhadze A.A. *Kvantovaya Elektron.*, **39**, 68 (2009) [*Quantum Electron.*, **39**, 68 (2009)].
- Mori W.B., Katsouleas T. *Phys. Rev. Lett.*, **69**, 3495 (1992).
- Korobkin V.V., Romanovskiy M.Yu., Trofimov V.A., Shiryaev O.B. *Kvantovaya Elektron.*, **43**, 232 (2013) [*Quantum Electron.*, **43**, 232 (2013)].
- Wang J.X., Scheid W., Hoelss M., Ho Y.K. *Phys. Rev. E*, **65**, 028501 (2002).
- McDonald K.T. *Phys. Rev. Lett.*, **80**, 1350 (1998).
- Mora P., Quesnel B. *Phys. Rev. Lett.*, **80**, 1351 (1998).
- Lin H., Wang C., Liu C.P. *Opt. Commun.*, **292**, 101 (2013).
- Esarey E., Sprangle P., Krall J., Ting A. *IEEE Trans. Plasma Sci.*, **24**, 252 (1996); Esarey E., Schroeder C.B., Leemans W.P. *Rev. Mod. Phys.*, **81**, 1229 (2009).
- Perry M.D., Mourou G. *Science*, **264**, 917 (1994); Mourou G., Barty C.P.J., Perry M.D. *Phys. Today*, **51**, 22 (1998).
- Lin H., Liu C.P., Du S.T., Wang C. *Appl. Phys. Res.*, **5** (3), 123 (2013).
- Rau B., Tajima T., Hojo H. *Phys. Rev. Lett.*, **78**, 3310 (1997).
- Sessler A.M. *Phys. Today*, **1**, 26 (1988).
- Ho Y.K., Feng L. *Phys. Lett. A*, **184**, 440 (1994).
- Ho Y.K., Wang J.X., Feng L., Scheid W., Hora H. *Phys. Lett. A*, **220**, 189 (1996).

34. Scheid W., Hora H. *Laser Part. Beams*, **7**, 315 (1989).
35. Goreslavsky S.P., Fedorov M.V., Kilpio A.A. *Laser Phys.*, **5**, 1020 (1995).
36. Galkin A.L., Korobkin V.V., Romanovskiy M.Yu., Trofimov V.A., Shiryaev O.B. *Phys. Plasmas*, **19**, 073102 (2012).
37. Korobkin V.V., Romanovskiy M.Yu., Trofimov V.A., Shiryaev O.B. *Laser Part. Beams*, **31**, 23 (2013).