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Autoresonance laser acceleration of guided "quasineutral" electron-positron beams

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A scheme for simultaneous laser acceleration of electrons and positrons to high energies is discussed. The acceleration is based on a cyclotron autoresonance between the particles and a linearly polarized radiation, propagating along an axial magnetic field. The nonlinear dynamics of an electron and a positron in the system is investigated for homogeneous fields and luminous radiation. The energy of the accelerated particles scales asymptotically as $0.82\{[B_z(100\ \text{kG})/\lambda(10\ \mu\text{m})]^{1/2}\alpha_x z(\text{m})\}^{2/3}$ GeV, where z and B_z are the distance and strength of the guide magnetic field and $\alpha_x = eE/(m\omega c)$ and λ are the normalized radiation field amplitude and its wavelength, respectively. It is shown that high-current "quasineutral" beams (plasmas) can be accelerated to high energies with low radiation losses.

I. INTRODUCTION

Acceleration of particles to high energies by high-power lasers has received considerable attention in recent years.[1–3] Several schemes for coupling between the laser field and relativistic charged beams have been proposed. Two major schemes applying the far field approach were discussed: the inverse free-electron laser[4] (IFEL) and the autoresonance laser accelerator[5] (ALA). In the IFEL accelerator the attainable energies are limited ($\lesssim 400$ GeV) due to radiation losses, and the accelerated beam current is limited as a result of transverse gradients of the wiggler field.

The ALA scheme is based on a cyclotron autoresonance between charged particles and an electromagnetic wave propagating along an axial guide magnetic field. This self-sustained resonance is achieved for luminous radiation and homogeneous fields. Both the nonlinear single-particle[5] and the collective[6] behavior of this system have been analyzed for the case of a circularly polarized laser radiation. It has been shown that the accelerated beam can be launched into the desired autoresonance regime through an appropriate transition region. The possibility of using the megagauss magnetic fields generated spontaneously in laser-produced plasmas for autoresonance acceleration was also suggested.[7] The ALA concept allows the acceleration of high-current beams to TeV energies with low radiation losses. However, generally speaking, the collective acceleration of dense beams is limited by space-charge effects. In this paper we propose a novel approach to cancel these effects by accelerating "quasineutral" beams (plasmas) to high energies applying the ALA concept. In particular, we consider the simultaneous acceleration of guided electrons and positrons by a linearly polarized electromagnetic radiation. Lately, the autoresonance acceleration of pairs by a nonpropagating electromagnetic wave has been discussed.[8] However, this approach cannot be extended to high energies because of synchrotron radiation losses.

The scope of the present work is as follows. The nonlinear electron dynamics in combined guide magnetostatic and linearly polarized laser radiation fields is considered analytically in Sec. II for homogeneous fields and luminous radiation. It is shown that at these conditions two autoresonances can be achieved. In Sec. III we analyze the symmetry between the autoresonant acceleration of a positron and an electron. This symmetry allows the simultaneous acceleration of electron-positron pairs to high energies. Finally, Sec. IV summarizes the results of this work.

II. THE ELECTRON DYNAMICS IN COMBINED GUIDE MAGNETOSTATIC AND LINEARLY POLARIZED LASER RADIATION FIELDS

Consider a plane electromagnetic wave propagating in the z direction along an axial homogeneous magnetostatic field $\mathbf{B}_z$. Let $\mathbf{k}_0$, ω_0' be the wave vector and the frequency of the wave, respectively. Define the following two sets of orthonormal vectors:

$$\left.\begin{aligned}\hat{e}_1 &= -\hat{e}_x\sin\phi + \hat{e}_y\cos\phi\\ \hat{e}_2 &= -\hat{e}_x\cos\phi - \hat{e}_y\sin\phi\\ \hat{e}_3 &= \hat{e}_z\end{aligned}\right\},\qquad \left.\begin{aligned}\hat{\hat{e}}_1 &= -\hat{e}_x\sin\phi - \hat{e}_y\cos\phi\\ \hat{\hat{e}}_2 &= -\hat{e}_x\cos\phi + \hat{e}_y\sin\phi\\ \hat{\hat{e}}_3 &= \hat{e}_z\end{aligned}\right\}, \tag{1}$$

where $\phi = k_0 z - \omega_0' t$. The pairs of vectors $\{\hat{e}_1,\hat{e}_2\}$ and $\{\hat{\hat{e}}_1,\hat{\hat{e}}_2\}$ are rotating in opposite directions as functions of

ϕ. The transformation between the two sets of vectors is

$$\hat{\hat{\mathbf{e}}}_1=-\cos(2\phi)\hat{\mathbf{e}}_1+\sin(2\phi)\hat{\mathbf{e}}_2\ ,$$
$$\hat{\hat{\mathbf{e}}}_2=\sin(2\phi)\hat{\mathbf{e}}_1+\cos(2\phi)\hat{\mathbf{e}}_2\ , \tag{2}$$
$$\hat{\hat{\mathbf{e}}}_3=\hat{\mathbf{e}}_3\ .$$

The electromagnetic fields in the systems can be written as

$$\mathbf{E}=-\frac{1}{c}\frac{\partial \mathbf{A}}{\partial t}\ ,\qquad \mathbf{B}=\nabla\times\mathbf{A}+B_z\hat{\mathbf{e}}_3\ , \tag{3}$$

where the vector potential is assumed to be linearly polarized in the x direction

$$\mathbf{A}=A_x(\cos\phi)\hat{\mathbf{e}}_x=-A\hat{\mathbf{e}}_2-A\hat{\hat{\mathbf{e}}}_2 \tag{4}$$

and $A=A_x/2$.

Thus the electromagnetic wave can be decomposed into right- and left-handed circularly polarized waves. The amplitude of each of these waves is half of the amplitude of the linearly polarized wave.

We shall discuss in the following the dynamics of a cold-electron beam in the described fields configuration, assuming that **E** and **B** are large enough and thus are not affected by the beam itself. The electrons momentum equation is

$$\left[\frac{\partial}{\partial t}+\mathbf{v}\cdot\nabla\right](m\gamma\mathbf{v})=-e\left[\mathbf{E}+\frac{\mathbf{v}}{c}\times\mathbf{B}\right]\ , \tag{5}$$

where $\gamma=(1-v^2/c^2)^{-1/2}$.

Equivalently, from Eq. (3),

$$\left[\frac{\partial}{\partial t}+\mathbf{u}\cdot\nabla\right](\gamma\mathbf{u}-\boldsymbol{\alpha})=\boldsymbol{\Omega}\times\mathbf{u}-(\nabla\boldsymbol{\alpha})\cdot\mathbf{u}\ , \tag{6}$$

where we use the following notations:

$$\boldsymbol{\Omega}=(eB_z/mc^2)\hat{\mathbf{e}}_3,\quad \boldsymbol{\alpha}=(e\mathbf{A}/mc^2),\quad \mathbf{u}=\mathbf{v}/c\ ,$$
$$\omega_0=\omega_0'/c,\quad \tau=ct\ .$$

According to Eqs. (2) and (4),

$$\boldsymbol{\alpha}=\alpha_x(\cos\phi)\hat{\mathbf{e}}_2=-\alpha\hat{\mathbf{e}}_2-\alpha\hat{\hat{\mathbf{e}}}_2=\alpha_1\hat{\mathbf{e}}_1+\alpha_2\hat{\mathbf{e}}_2\ , \tag{7}$$

where

$$\alpha_1=-\alpha\ \sin(2\phi),\quad \alpha_2=-2\alpha\cos^2\phi\ ,$$
$$\alpha=eA_x/2mc^2=\alpha_x/2\ .$$

Next we write Eq. (6) in components in the frame $\{\hat{\mathbf{e}}_j\}$,

$$\frac{d}{d\tau}(\gamma u_1-\alpha_1)+(\omega_0-k_0u_3)(\gamma u_2-\alpha_2)=-\Omega u_2\ , \tag{8}$$

$$\frac{d}{d\tau}(\gamma u_2-\alpha_2)-(\omega_0-k_0u_3)(\gamma u_1-\alpha_1)=\Omega u_1\ , \tag{9}$$

$$\frac{d}{d\tau}(\gamma u_3)=k_0(\alpha_2u_1-\alpha_1u_2)\ . \tag{10}$$

The energy of the electron evolves according to the equation

$$\frac{d\gamma}{d\tau}=u_1\frac{\partial\alpha_1}{\partial\tau}+u_2\frac{\partial\alpha_2}{\partial\tau}+\omega_0(\alpha_2u_1-\alpha_1u_2)\ . \tag{11}$$

For homogeneous and stationary fields ($\alpha,\Omega=$const), Eqs. (8)–(11) can be simplified by using Eq. (7),

$$(\gamma u_1)^{\cdot}=k_0(u_3-u_p)(\gamma u_2+2\alpha\sin^2\phi)-\Omega u_2\ , \tag{12}$$

$$(\gamma u_2)^{\cdot}=-k_0(u_3-u_p)[\gamma u_1-\alpha\sin(2\phi)]+\Omega u_1\ , \tag{13}$$

$$(\gamma u_3)^{\cdot}=(\dot\gamma/u_p)=-2k_0\alpha(u_1\sin\phi+u_2\cos\phi)\sin\phi\ , \tag{14}$$

where $(\ )^{\cdot}\equiv d(\ )/d\tau$. These are the desired equations describing the electron dynamics in the system. Equation (14) yields the following constant of motion:

$$L_0=\gamma(1-u_pu_3)=\text{const}\ . \tag{15}$$

Note that Eqs. (12)–(14) include expressions containing ϕ. Therefore, their direct time integration is difficult since ϕ is a function of z and

$$z(\tau)=\int_0^\tau u_3d\tau\ . \tag{16}$$

One can overcome this difficulty by replacing the time derivatives in these equations by phase derivatives. For an arbitrary variable a, we have

$$(a)^{\cdot}=\dot\phi(da/d\phi)\equiv\dot\phi(a)_\phi=k_0(u_3-u_p)(a)_\phi\ . \tag{17}$$

Therefore, Eqs. (12) and (13) can be rewritten as

$$(\gamma u_1)_\phi=l(\gamma u_2)+2\alpha\sin^2\phi\ , \tag{18}$$

$$(\gamma u_2)_\phi=-l(\gamma u_1)+\alpha\sin(2\phi)\ , \tag{19}$$

where $l=1+\Omega/[k_0\gamma(u_p-u_3)]$. Equations (18) and (19) exhibit simple resonance behavior when $l=$const. According to Eq. (15), this takes place only for luminous radiation, i.e., when $u_p=1$. We will consider the electron dynamics in this case in the following. Equations (18) and (19) yield a "forced harmonic oscillator" type of differential equation for (γu_1)

$$(\gamma u_1)_{\phi\phi}+l^2(\gamma u_1)=(2+l)\alpha\sin(2\phi)\ , \tag{20}$$

where $l=1+\Omega/(k_0L_0)$. The last equation shows the existence of two self-sustained cyclotron resonances in the system, namely,

$$l=0\leftrightarrow\Omega=-k_0L_0,\quad l=2\leftrightarrow\Omega=k_0L_0\ . \tag{21}$$

The two autoresonances differ in the direction of the magnetic field. First, we shall discuss off-resonance conditions. The general solution of Eq. (20) for $l\neq0,2$ is

$$\gamma u_1=\frac{\alpha}{(l-2)}\sin(2\phi)+\gamma_0u_{10}\cos(l\phi)+\left[\gamma_0u_{20}-\frac{2\alpha}{l(l-2)}\right]\sin(l\phi)\ , \tag{22}$$

where the initial conditions are $\gamma u_j(\phi=0)=\gamma_0u_{j0}$ ($j=1,2$) and according to Eq. (18), $(\gamma u_1)_\phi=l\gamma_0u_{20}$ at $\phi=0$. By substituting Eq. (22) into Eq. (19) we get

$$\gamma u_2=\frac{\alpha}{(l-2)}[\cos(2\phi)-1]-\gamma_0 u_{10}\sin(l\phi)+\gamma_0 u_{20}\cos(l\phi)-\frac{2\alpha}{l(l-2)}[\cos(l\phi)-1]\ . \quad (23)$$

According to Eq. (15),

$$u_3=1-L_0/\gamma\ , \quad (24)$$

and since $\gamma=(1-u_1^2-u_2^2-u_3^2)^{-1/2}$, we obtain

$$\gamma=\frac{1}{2L_0}[1+(\gamma u_1)^2+(\gamma u_2)^2+L_0^2]\ . \quad (25)$$

Thus, it is evident that $\gamma\to\infty$ if $(\gamma u_1)\to\infty$ as $|\phi|\to\infty$. The function $\gamma=\gamma(\phi)$ for $l\neq 0,2$ can be found from Eqs. (22) and (25). The energy of the electron is limited in this case since the solutions for $(\gamma u_1),(\gamma u_2)$ are bounded. The maximum value of γ can be found by setting $\dot{\gamma}=0$ in Eq. (14),

$$\gamma_{\max}=\gamma(\phi^*),\quad \phi^*=\arctan(-u_2/u_1)\ . \quad (26)$$

For a given l, the maximum energy gain by the electron can be found from Eqs. (22), (23), (25), and (26).

Next we consider the electron dynamics at the autoresonance $l=0$. In this case the solutions of Eqs. (18) and (19) are

$$\gamma u_1=\gamma_0 u_{10}+\alpha[\phi-\tfrac{1}{2}\sin(2\phi)]\ , \quad (27)$$

$$\gamma u_2=\gamma_0 u_{20}+\alpha\sin^2\phi\ . \quad (28)$$

According to Eq. (1), the transverse velocity components in the laboratory frame are

$$\gamma u_x=\gamma_0 u_{x0}-\alpha\phi\sin\phi\ ;$$

$$\gamma u_y=\gamma_0 u_{y0}-\alpha(\sin\phi-\phi\cos\phi)\ .$$

In order to find the time dependence of the electron's energy we use the equalities

$$\gamma\dot{\gamma}=k_0\gamma(u_3-1)\gamma_\phi=-k_0L_0\gamma_\phi \quad (29)$$

and, from Eqs. (14) and (27)–(29),

$$\gamma_\phi=(2\alpha\sin\phi/L_0)[\gamma_0(u_{10}\sin\phi+u_{20}\cos\phi)+\alpha\phi\sin\phi]\ . \quad (30)$$

The integration of the last equation yields

$$\gamma=\gamma_0+(\alpha/2L_0)\{-\gamma_0[u_{10}\sin(2\phi)+u_{20}\cos(2\phi)]+2\gamma_0u_{10}\phi+\gamma_0u_{20}+\alpha[\phi^2-\phi\sin(2\phi)+\sin^2\phi]\}\ . \quad (31)$$

The time dependence of ϕ for $l=0$ can be found from

$$\dot{\phi}=k_0(u_3-1)=\Omega/\gamma\ , \quad (32)$$

yielding

$$\tau=\frac{1}{\Omega}\int\gamma\dot{\phi}d\tau=\frac{1}{\Omega}\int\gamma(\phi)d\phi\ . \quad (33)$$

By substituting Eq. (31) into Eq. (33) we have

$$\tau=\frac{\gamma_0}{\Omega}\phi+\left[\frac{k_0\alpha}{2\Omega^2}\right]\left[\frac{\gamma_0}{2}[u_{10}\cos(2\phi)-u_{20}\sin(2\phi)]-\gamma_0u_{10}\phi^2-\gamma_0u_{20}\phi-\tfrac{1}{2}\gamma_0u_{10}-\alpha\left[\frac{\phi^3}{3}+\tfrac{1}{2}\phi[\cos(2\phi)+1]-\tfrac{1}{2}\sin(2\phi)\right]\right]\ . \quad (34)$$

Equations (24), (27), (28), (31), and (34) describe the nonlinear dynamics of the electron at the resonance $l=0$. Asymptotically, for $\phi\to-\infty$, one obtains

$$\mathbf{u}_\perp\simeq u_1\hat{\mathbf{e}}_1\simeq\frac{\alpha\phi}{\gamma}\hat{\mathbf{e}}_1\simeq-\left[\frac{2L_0}{\gamma}\right]^{1/2}\hat{\mathbf{e}}_1\ , \quad (35)$$

$$\phi\simeq-\left[\frac{6\Omega^2}{k_0\alpha^2}\tau\right]^{1/3}=-7.22\left[\left[\frac{B_z(100\ \text{kG})}{\alpha}\right]^2\lambda_0(10\ \mu\text{m})\tau(\text{m})\right]^{1/3}, \quad (36)$$

$$\gamma\simeq\left[\frac{3|\alpha|}{\sqrt{2}}(k_0|\Omega|)^{1/2}\tau\right]^{2/3}=2.6\times10^3\left[|\alpha|\left|\frac{B_z(100\ \text{kG})}{\lambda_0(10\ \mu\text{m})}\right|^{1/2}\tau(\text{m})\right]^{2/3}, \quad (37)$$

where $\lambda_0=(2\pi/k_0)$ is the radiation wavelength.

These asymptotic results coincide with the results obtained previously[5] for circularly polarized radiation $\boldsymbol{\alpha}=\alpha_2\hat{\mathbf{e}}_2$, if we perform the transformation $-\alpha=-(\alpha_x/2)\to\alpha_2$. At the resonance $l=0$ the electron is accelerated by the circularly polarized component $-\alpha\hat{\mathbf{e}}_2$ of $\boldsymbol{\alpha}$ [see Eq. (7)]. Therefore the accelerating amplitude is only one-half of the amplitude of the linearly polarized radiation.

Next, we consider the autoresonance $l=2$. For this case, Eq. (20) has the solution

$$\gamma u_1=(\gamma_0u_{20}+\alpha/2)\sin(2\phi)+(\gamma_0u_{10}-\alpha\phi)\cos(2\phi)\ , \quad (38)$$

and using Eq. (19) we get

$$\gamma u_2=\gamma_0u_{20}\cos(2\phi)-(\gamma_0u_{10}-\alpha\phi)\cos(2\phi)\ . \quad (39)$$

It is convenient now to make the transformation to the frame $\{\hat{\hat{\mathbf{e}}}_j\}$ by defining $\mathbf{u}_\perp=\tilde{u}_1\hat{\hat{\mathbf{e}}}_1+\tilde{u}_2\hat{\hat{\mathbf{e}}}_2$, where

$$\tilde{u}_1=u_1\hat{\hat{\mathbf{e}}}_1\cdot\hat{\mathbf{e}}_1+u_2\hat{\hat{\mathbf{e}}}_1\cdot\hat{\mathbf{e}}_2=-u_1\cos(2\phi)+u_2\sin(2\phi)\ , \quad (40)$$

$$\tilde{u}_2=u_1\hat{\hat{\mathbf{e}}}_2\cdot\hat{\mathbf{e}}_1+u_2\hat{\hat{\mathbf{e}}}_2\cdot\hat{\mathbf{e}}_2=u_1\sin(2\phi)+u_2\cos(2\phi)\ , \quad (41)$$

and, according to Eqs. (38) and (39),

$$\gamma\tilde{u}_1=-\gamma_0u_{10}+\alpha[\phi-\tfrac{1}{2}\sin(2\phi)]\ , \quad (42)$$

$$\gamma\tilde{u}_2=\gamma_0u_{20}+\alpha\sin^2\phi\ . \quad (43)$$

Therefore, for $l=2$ and $u_1(\phi=0)=u_{10}$ the electron transverse momentum components in the frame $\{\hat{\hat{e}}_j\}$ are identical with those for $l=0$ and $u_1(\phi=0)=-u_{10}$ in the frame $\{\hat{e}_j\}$. Furthermore, Eqs. (14), (38), and (39) yield

$$\gamma_\phi=(2\alpha\sin\phi/L_0)[\gamma_0(u_{20}\cos\phi-u_{10}\sin\phi)+\alpha\phi\sin\phi] \ . \quad (44)$$

Thus, from Eqs. (30) and (44) we conclude that the autoresonance acceleration at $l=0$ and $u_1(\phi=0)=u_{10}$ is the same as the autoresonance acceleration at $l=2$ and $u_1(\phi=0)=-u_{10}$. In both cases the time-dependences of γ, u_3, and ϕ are identical. Nevertheless, since the direction of the magnetic field is reversed, the electron rotates in opposite directions around the z axis in these resonances. Asymptotically, for $l=2$,

$$\gamma\mathbf{u}_\perp \simeq \frac{\alpha\phi}{\gamma}\hat{\hat{e}}_1 \simeq -\left[\frac{2L_0}{\gamma}\right]^{1/2}\hat{\hat{e}}_1 \ . \quad (45)$$

In this case the electron is accelerated by the circularly polarized component $-\alpha\hat{\hat{e}}_2$ of $\boldsymbol{\alpha}$. The accelerating electric fields corresponding to the vector-potential components $(-\alpha\hat{e}_2),(-\alpha\hat{\hat{e}}_2)$ rotate in the directions $\hat{e}_1$ and $\hat{\hat{e}}_1$, respectively. Asymptotically, as follows from Eqs. (35) and (45), the electron rotates in resonance with the electric field associated with the appropriate component of $\boldsymbol{\alpha}$, according to the direction of $\boldsymbol{\Omega}$. Therefore, asymptotically $(|\phi|\gg 1)$ the acceleration is the same for a linearly polarized radiation with a certain amplitude and a circularly polarized radiation with half of this amplitude. Thus in accordance with previous calculations,[5] the radiation losses in a 1 TeV autoresonance laser accelerator can be negligible. Finally, Fig. 1 shows an example of an autoresonance acceleration by a Nd:glass laser with $\alpha_x=2$, $(\alpha=1)$, along a 100-kG magnetic field [see Eqs. (31) and (34)]. The z dependence of γ and ϕ is shown for an electron beam with $\mathbf{u}_{\perp 0}=0$.

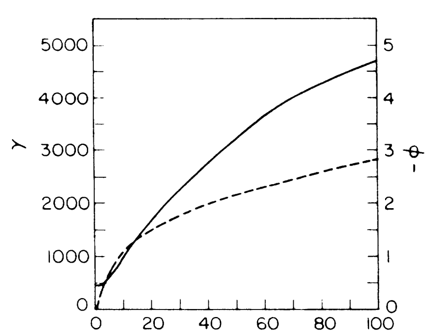

FIG. 1. Autoresonance acceleration of an electron beam with $\gamma_0=470$, $u_{10}=u_{20}=0$, by a linearly polarized Nd:glass radiation (with wavelength 1.06 μm) with $\alpha=1$ (intensity of $\sim 10^{18}$W/cm^2) along a 100-kG magnetic field. The z dependence of γ (solid line) and ϕ (dashed line) is presented.

III. SIMULTANEOUS ACCELERATION OF ELECTRONS AND POSITRONS

The dynamics of a positron in the system is described by Eqs. (8) and (11) with reversed signs for α_1, α_2, and Ω. For luminous radiation and homogeneous fields, the positron obeys

$$(\gamma u_1)_{\phi\phi}+l_p^2(\gamma u_1)=-(2+l_p)\alpha\sin(2\phi) \ , \quad (46)$$

where $l_p=1-\Omega/(k_0L_0)$. Therefore, similarly to an electron, the positron has resonances at $\Omega=\pm k_0L_0$. After replacing l by l_p and changing the sign of α, Eqs. (22), (23), (31), (34), (38), and (39) describe the nonlinear dynamics of the positron at the resonances $l_p=0,2$. At an axial magnetic field $\Omega=k_0L_0$, the electron and the positron are at the resonances $l=2$ and $l_p=0$, respectively, whereas at $\Omega=-k_0L_0$, the resonances are $l=0$ and $l_p=2$. For a given resonant magnetic field the two particles are rotating in opposite directions and are accelerated in phase with different circularly polarized components of $\boldsymbol{\alpha}$. Therefore, electrons and positrons can be accelerated simultaneously to high energies. At resonance the evolution of γ, u_3, and ϕ will be the same for electrons and positrons with zero initial transverse velocity $(u_{10}=u_{20}=0)$. At these conditions, if initially the density of the electrons N_{e^-} $(z=0)$ is equal to the density of the positrons N_{e^+} $(z=0)$ in a beam, then,

$$N_{e^+}(z)=N_{e^-}(z) \quad (47)$$

for each $z>0$. In conclusion, by using a linearly polarized laser radiation, one can accelerate quasineutral beams (plasmas) containing equal densities of electrons and positrons. Figure 2 presents a possibility for collinear autoresonance acceleration of quasineutral electron-positron beams. The beams propagating in the $-z$ and the $+z$ directions are accelerated by the $\{l=2,l_p=0\}$ and the $\{l=0,l_p=2\}$ resonances, respectively.

IV. DISCUSSION AND CONCLUSIONS

In this work we have considered the nonlinear electron (positron) dynamics in cyclotron autoresonance with a

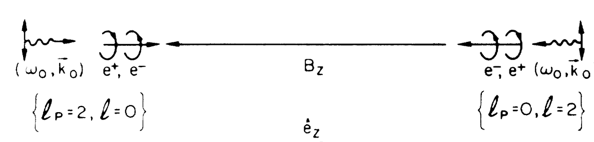

FIG. 2. Autoresonance laser acceleration of oppositely propagating quasineutral electron-positron beams along an axial magnetic field B_z. The types of resonances between the particles and the linearly polarized lasers are indicated.

linearly polarized laser radiation propagating along an axial magnetic field. There are two resonances for the electron (positron) in the system $l=2,0(l_p=0,2)$, which differ in the direction of the magnetic field. Accordingly, the particle is accelerated in resonance with one of the circularly polarized components of the laser radiation. Therefore, the acceleration by a linear polarized radiation of a given amplitude coincides asymptotically with the acceleration by a circularly polarized radiation with half of this amplitude [see Eqs. (35)—(37) and Ref. 5]. It has been shown that simultaneous autoresonance acceleration of electrons and positrons to high energies can be achieved. This result cannot be obtained within the existing schemes for laser acceleration of particles.[1,2] In other schemes, such as the plasma accelerators[9–11] or the inverse free electron laser,[4] a positron is deccelerated in a phase at which an electron is accelerated. In contrast, the autoresonance laser acceleration scheme allows continuous acceleration of quasineutral beams to high energies with low radiation losses. In particular, these beams can contain equal densities of electrons and positrons or other plasmas with oppositely charged particles of equal $|q|/m$ (e.g., positive and negative ions). Such beams can be accelerated without the usual space-charge divergence effects. The acceleration of high-current beams, thus, seems to be feasible. Finally, we note that the autoresonance acceleration mechanism might play an important role in cosmic-ray acceleration in astrophysical objects, where magnetostatic fields and intense electromagnetic radiation coexist.

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