Particle manipulation with nonadiabatic ponderomotive forces[a)]

I. Y. Dodin[b)] and N. J. Fisch
Department of Astrophysical Sciences, Princeton University, Princeton, New Jersey 08544

(Received 3 November 2006; accepted 16 November 2006; published online 21 March 2007)

Average, or ponderomotive potentials effectively seen by particles in oscillating fields allow advanced techniques of particle manipulation inaccessible with static potentials. In strongly inhomogeneous fields the ponderomotive force is phase dependent, and the particle dynamics resembles that of a quantum object in a conservative barrier. Probabilistic transmission through a ponderomotive potential is then possible and can be used for particle beam slicing. Resonant fields can also cool and trap particles exhibiting natural oscillations (e.g., Larmor rotation), as well as transmit them asymmetrically; hence, acting as one-way walls. An approximate integral of particle motion is found for this case and a new ponderomotive potential is introduced accordingly. *© 2007 American Institute of Physics*. [DOI: 10.1063/1.2436149]

I. INTRODUCTION

The 1997 Nobel Prize in physics was awarded for the invention of methods of atom manipulation by laser light.[1] Similar capabilities also apply to other objects, ranging from molecules to micrometer-sized particles, and permit one to selectively and stably trap particles, levitate them against gravity, channel particles along laser beams, and use them as sensitive probes for measuring optical, electric, magnetic, viscous drag, and gravity forces.[2] The new tools of particle control yield present and potential applications in a variety of subjects such as light scattering, cloud physics, quantum optics, isotope separation, and high-resolution spectroscopy.[3] What we offer here is the extension of these techniques based on classical interactions of undamped particles with intense electromagnetic fields.[4]

Even without a bias, an ac field can exert a nonzero time-averaged force on a particle.[5–7] This so-called ponderomotive force consists of two components: the dipole force due to the inhomogeneity of the field envelope and the light pressure due to the radiation scattering off the particle. Often, the light pressure is negligible (the particle is "undamped"), and the induced dipole moment of the particle $\mathbf{p}$ is nearly a local, or "adiabatic" function of the field $\mathbf{E}$. In the linear approximation, the two conditions yield $\mathbf{p}=\boldsymbol{\alpha}\cdot\mathbf{E}$, where the polarizability tensor $\boldsymbol{\alpha}$ is Hermitian. The force on the particle can then be described in terms of the ponderomotive potential Φ, equal to the average energy of the dipole-field interaction:[8]

$$\Phi = -\frac{1}{4}(\mathbf{E}^* \cdot \boldsymbol{\alpha} \cdot \mathbf{E}). \tag{1}$$

Below we assume a stationary field ($\mathbf{k}=0$, $\partial_t|E|\equiv 0$); thus, Φ is a function of space only. For example, in the absence of additional forces, an elementary particle with charge e and mass m has $\boldsymbol{\alpha}=-e^2/m\omega^2\mathbf{I}$, so one recovers the well-known formula $\Phi=e^2|E|^2/4m\omega^2$.[5,9]

Under the above assumptions, ac fields act like static potentials, which yields numerous applications including atomic traps,[1] rf plugs,[6,10] low-frequency mode stabilization,[11] and edge control in fusion plasmas.[12] Violating the approximation (1), though, renders extra flexibility, allowing for more advanced and otherwise inaccessible techniques of particle manipulation. References 1 illustrate how the additional capabilities result from dissipation. What we contemplate here is how similar effects can be practiced on undamped particles—via breaking the adiabatic approximation.

To offer new tools for manipulating classical particles with ac fields and advance the analytical treatment of nonadiabatic ponderomotive forces is the purpose of this paper. The work represents a review of our selected results reported in Refs. 13–22, which we re-examine here from a unifying standpoint. We show that in strongly inhomogeneous fields, the ponderomotive force is phase dependent, and the particle dynamics resembles that of a quantum object in a conservative barrier. Probabilistic transmission through a ponderomotive potential is then possible and can be used for particle beam slicing. We demonstrate that resonant fields can also cool and trap particles exhibiting natural oscillations (e.g., Larmor rotation), as well as transmit them asymmetrically, hence acting as one-way walls. We find an approximate integral of particle motion for this case and introduce a new ponderomotive potential accordingly.

The paper is organized as follows. In Sec. II, we explain the analogy between the nonadiabatic and quantum dynamics. In Sec. III, we discuss the particle beam slicing and nonadiabatic tunneling of particles through ponderomotive barriers. In Sec. IV, we offer an analytical model to describe particle dynamics in resonant barriers. In Sec. V, we discuss cooling and trapping of particles by ponderomotive forces. In Sec. VI, we explain the operation of ponderomotive one-way walls and the current drive effect. In Sec. VII, we summarize our results.

[a)]Paper WI1 2, Bull. Am. Phys. Soc. **51**, 332 (2006).
[b)]Invited speaker.

II. QUANTUM ANALOGY

A multiscale approach to describing the particle dynamics driven by intense electromagnetic radiation relies on separating fast oscillatory motion of the particle from its slow translational motion. Assume that the parameters of the oscillations vary along the particle trajectory $\mathbf{s}(t)$ on a sufficiently large scale L; that is,

$$\epsilon \equiv \max\{\lambda/L, d\lambda/ds\} \ll 1, \tag{2}$$

where λ is the particle average displacement on the oscillation period. The translational motion is then conveniently described in terms of the so-called "guiding center" variables, for which the explicit time dependence is removed from the motion equations to any power in ϵ.[13,23] This technique is widely used in theoretical and computational plasma physics to study particle dynamics in rf and laser fields;[24] what is in practice often missed, though, is the intrinsically limited accuracy of the guiding center approximation. The transformation to the new coordinates is an asymptotic procedure with an exponentially small error in ϵ.[13,20] The true drift coordinates are then in principle definable only with limited accuracy; thus, the dynamics of the guiding center as a quasiparticle may not follow the laws of classical mechanics. Strikingly, what this dynamics resembles instead is the motion of a quantum object. For the particle motion in a nonuniform magnetic field, a similar analogy was previously drawn by Varma, as reviewed in Ref. 25. However, the explanation of the "macroquantum" effects in terms of the actual wave functions remains controversial.[26] In contrast to Varma's quantum approach, we show that purely classical particles can exhibit *quantum-like* effects.

A vivid manifestation of quantum-like dynamics is the existence of quantized eigenstates for particle bounce oscillations in a confining ponderomotive potential, which can be explained as follows. At finite ϵ, each time a particle bounces off a ponderomotive wall, it either gains or loses energy, depending on its velocity and phase. At sufficiently large ϵ, stochastic oscillations may develop, in which case the particle would be heated by the ac field until it leaves the interaction region. Inside the stochasticity domain, though, there exists a countable set of closed phase space trajectories, for which the amount of energy gained by a particle per bounce period equals the energy transferred back to the field. As ϵ increases, stability is preserved in the first place *near* these periodic orbits[27] (Fig. 1). Hence, rather than destroying confinement, the resonant interaction actually enhances it selectively in particular phase space regions. In a perturbed system, the confinement time will vary drastically depending on where the particle is located among those regions; this might explain similar effects reported in Ref. 25.

Resonant orbits are somewhat like the stationary eigenstates of a quantum particle in a conservative potential, which similarly exhibits discrete energy spectrum in a steady state. In Ref. 15, we show that the necessary condition for these orbits in the nearly adiabatic ("quasiclassical") domain follows the Bohr-Sommerfeld rule

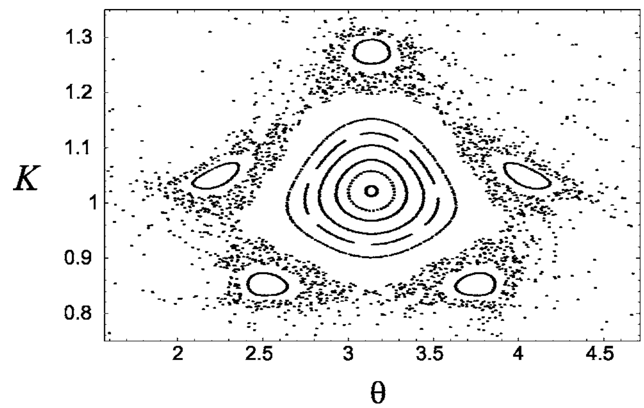

FIG. 1. Poincaré mapping (fragment) for the particle motion governed by $m\ddot{z}=eE(z)\sin\omega t$, with $E(z)=E_0\sinh^2(z/L)$, $\hat{\epsilon}\equiv eE/m\omega^2L=2$. Here, $K \equiv m\dot{z}^2/2$ a.u. and $\theta\equiv\omega t$ are taken at particle crossing $z=0$ with $\dot{z}>0$. Stable bounce oscillations (continuous curves) are observed near periodic, or resonant orbits (one in the center and five on the sides) considered as "eigenstates;" chaos is developed at the periphery. At larger $\hat{\epsilon}$, only particles near the central orbit are confined.

$$\oint k dz \approx 2\pi n \tag{3}$$

(here, $k=2\pi/\lambda$, and n is an integer), due to the reason akin to that for a truly quantum particle. Like a quantum object, the particle guiding center is not a zero-dimensional entity but is assigned a phase, which is the phase of the particle oscillations in the ac field. The distance λ, which determines the "uncertainty" of the guiding center location, can be naturally treated as the effective de Broglie wavelength. One may say then that the ponderomotive potential approximation holds when the guiding center coordinates are well defined.

The quantum analogy extends further and applies also to free (nonconfined) particles. The average force on the guiding center is proportional to the amplitude of the particle induced oscillations. Depending on the initial phase, a fast enough particle ($\epsilon \gtrsim 1$) may not have sufficient time to gain quiver energy from the field, in which case, neither will it experience significant average acceleration. Such a particle will then be able to penetrate through "classically forbidden" regions, just like a quantum particle having a de Broglie wavelength of the order of the field scale. In Sec. III, we consider this nonadiabatic "tunneling" in more detail. In Sec. V, we also discuss a related quantum-like effect, such as particle reflection from an attractive ponderomotive potential.

III. NONADIABATIC TUNNELING

Consider nonadiabatic tunneling for elementary particles, for simplicity assuming a localized one-dimensional ac envelope $E(z)$ with no other fields imposed. In this case $\lambda=2\pi v/\omega$, where v is the particle drift velocity; hence, the adiabaticity parameter ϵ coincides with $\varsigma\equiv 2\pi v/\omega L$ being the ratio of the particle displacement on the *field* period to the field scale L.[28] A ponderomotive barrier with $\varsigma\to 0$ acts just like a static potential, so the "quasienergy"

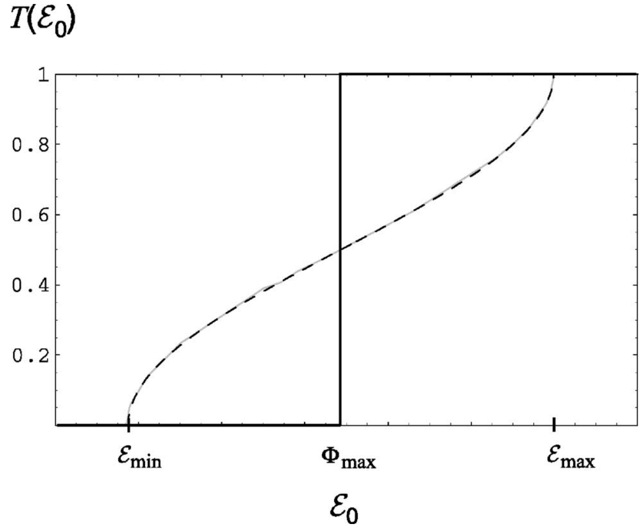

FIG. 2. Transmission coefficient $T(\mathcal{E}_0)$ for a particle scattering off a ponderomotive barrier with $E(z)=E_0\exp(-z^2/L^2)$, $eE/m\omega^2L=0.1$: numerical (solid gray) and analytical [Eq. (5), using a fitting parameter $\varepsilon\approx 5.1\times 10^{-5}\Phi_{\max}$] results (dashed). Also shown is the "adiabatic" step-function approximation (solid black). $\Phi_{\max}$ is calculated according to Eq. (1), with higher-order corrections taken into account (Ref. 13).

$$\mathfrak{E}=\mathcal{E}+\Phi \tag{4}$$

is conserved, where $\mathcal{E}=mv^{2/2}$. (In Ref. 13, we show that $\mathfrak{E}$ can be considered as an adiabatic invariant of particle motion.) The barrier then transmits all particles with initial energy $\mathcal{E}_0>\Phi_{\max}$ and reflects those with $\mathcal{E}_0<\Phi_{\max}$, resulting in a step-like transmission coefficient $T(\mathcal{E}_0)=\Theta(\mathcal{E}_0-\Phi_{\max})$. As ς is increased, though, both energy and phase determine the transmission, so the phase-averaged $T(\mathcal{E}_0)$ becomes a continuous function.

An analogy with quantum tunneling through a static potential can be drawn in this case (Sec. II); yet $T(\mathcal{E}_0)$ is of algebraic form here rather than exponential. For a single-humped $E(z)$ with small ς, the transmission coefficient is derived analytically[13] and reads

$$T\approx\frac{1}{2}-\frac{1}{\pi}\arcsin\frac{\Phi_{\max}-\mathcal{E}_0}{\varepsilon}, \tag{5}$$

where ε is a functional of $E(z)$, such that $\varepsilon(\varsigma\to 0)=0$.[8] At ς approaching zero, Eq. (5) yields the adiabatic formula $T(\mathcal{E}_0)=\Theta(\mathcal{E}_0-\Phi_{\max})$. At nonzero ς, one has $T\propto\sqrt{\mathcal{E}_0-\mathcal{E}_{\min}}$ for $\mathcal{E}_0$ slightly above $\mathcal{E}_{\min}\equiv\Phi_{\max}-\varepsilon$, and $T\propto\sqrt{\mathcal{E}_{\max}-\mathcal{E}_0}$ for $\mathcal{E}_0$ slightly below $\mathcal{E}_{\max}\equiv\Phi_{\max}+\varepsilon$. Transmission is impossible at $\mathcal{E}_0<\mathcal{E}_{\min}$; reflection is impossible at $\mathcal{E}_0>\mathcal{E}_{\max}$ (Fig. 2).

Since particle transmission is phase-dependent, a uniform monoenergetic beam sent toward a barrier, say, from the left will be "sliced" by the ac field, and periodic particle bunches will appear on the right. Spatial spreading of these bunches, which occurs due to the phase modulation, can be suppressed if ultrarelativistic energies are employed ($\gamma_0\equiv\mathcal{E}_0/mc^2\gg 1$). Accounting for the relativistic corrections,[21,22] the normalized field amplitude $a=eE/mc\omega$ required to produce nonadiabatic tunneling is estimated as $a\sim\gamma_0$. For electrons with energies of several MeV,[29] appropriate laser fields with the wavelength of about one micron are available in laboratory.[30] Hence, subfemtosecond or even attosecond electron bunches can be produced with the existing technology using nonadiabatic ponderomotive barriers. The suggested method supplements related techniques employing intense laser-plasma interactions[31] and allows similar applications in generating extremely collimated x-ray pulses.

IV. EFFECTIVE POTENTIAL FOR A RESONANT FIELD

The particle motion in an ac field can be nonadiabatic also at $\varsigma\to 0$, if the particle exhibits natural oscillations at frequency Ω close to ω. (Examples of such oscillations are Larmor rotation in a dc magnetic field, quantum intra-atomic oscillations, molecular vibrations, and others.) In this case, the field can gradually pump up the particle internal degree of freedom, so the approximation of a local response does not hold. In other words, the adiabaticity parameter (2) is not small near the resonance, for here it must be defined with $\lambda=v\tau_{\mathrm{b}}$ being the average displacement on the *beat* period $\tau_{\mathrm{b}}=2\pi/|\omega-\Omega|$; hence, Eq. (1) becomes inapplicable.

Surprisingly, the corresponding dynamics is tractable analytically. In Ref. 14 we show that, given $\varsigma\ll 1$, a particle exhibiting linear oscillations of an arbitrary nature "sees" the effective potential

$$\Phi_{\mathrm{eff}}=\Phi-J\Delta\omega=\frac{\Delta\omega}{\Omega}(\mathcal{W}_{\mathrm{a}}-\mathcal{W}). \tag{6}$$

Here, $J=\mathcal{W}/\Omega$ is the action of the particle natural oscillations; $\mathcal{W}$ is the energy of these oscillations; $\Delta\omega=\omega-\Omega$ is the beat frequency; and $\mathcal{W}_a\propto 1/\Delta\omega^2>0$ is the energy of the formally introduced adiabatic oscillations linear to $\mathbf{E}$. In Ref. 14, we also generalize Eq. (6) to the case when more than one eigenfrequency Ω is present.

At nonadiabatic interaction ($\epsilon\gtrsim 1$), the variations of J, which would otherwise be conserved,[14,20] are nonlocal. Therefore, Φ_{eff} cannot be expressed as a single-valued function of space and hence is not quite a potential in the ordinary sense. Yet the quasi-energy

$$\mathfrak{E}=\mathcal{E}+\Phi_{\mathrm{eff}} \tag{7}$$

is an adiabatic invariant here, as also follows from the conservation of the total number of quanta and the energy in the "particle-field" system.[14,18,32] Contrary to the case in which no internal oscillations are present (Sec. III), now the conservation of $\mathfrak{E}$ by itself is insufficient for calculating particle trajectories. Nonetheless, Eq. (7) is enough to predict the general features of the particle energy exchange with the ac field. Based on those, applications of practical interest are proposed in Secs. V and VI.

V. COOLING AND TRAPPING

Breaking the adiabaticity allows irreversible energy exchange between particles and the ac field. If the radiation is redshifted from the resonance frequency Ω, thermal particles lose their drift energy as they scatter off a nonadiabatic ponderomotive barrier, regardless of their actual trajectories.

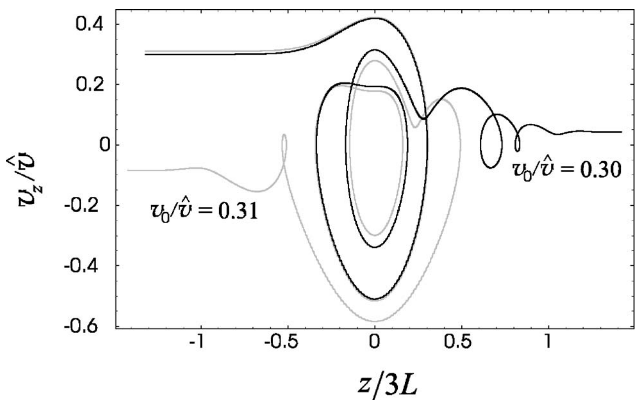

FIG. 3. Longitudinal velocity v_z vs z for a particle being trapped and released by an attractive ponderomotive barrier with $\mathbf{E}(z)=\hat{\mathbf{x}}E_0\exp(-z^2/2L^2)$ in a uniform dc magnetic field $\mathbf{B}_0=\hat{\mathbf{z}}B_0$ with $eB_0/mc\equiv\Omega>\omega$: $v_z=0.30\hat{v}$ (black) and $v_z=0.31\hat{v}$ (gray). Here, $\hat{v}\equiv(e|E_0|/m\omega)\sqrt{\Lambda}\sim(|\Phi|_{\max}/m)^{1/2}$, $eE_0/mc\omega=10^{-3}$, $\Lambda\equiv|1-\Omega/\omega|^{-1}=100$, $\hat{\epsilon}\equiv\hat{\lambda}/L=3$, $\hat{\lambda}=\hat{v}/|\omega-\Omega|$.

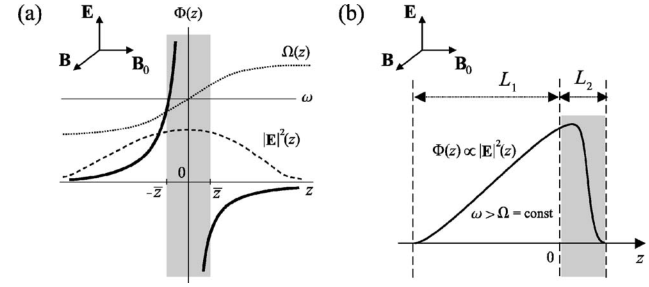

FIG. 4. Ponderomotive one-way walls. (a) Asymmetry is due to the inhomogeneity of the dc magnetic field $\mathbf{B}_0$. (b) Asymmetry is due to the inhomogeneity of the ac field $\mathbf{E}$. Here, $\Phi=e^2|E|^2/4m(\omega^2-\Omega^2)$, assuming linear polarization (Refs. 18 and 20); $\Omega=eB_0/mc$, $L_1\gg\lambda\gg L_2$, $\lambda=v_z/|\omega-\Omega|$. Shaded are regions of nonadiabatic interaction, i.e., $\bar{z}=\sqrt{Lv_z/\omega}$, according to Eq. (2). The effective heights of the barrier $\Phi_{\max}$ for particles incident from the left are: (a) $\Phi_{\max}\sim\Phi(|\bar{z}|)$,[18] and (b) $\Phi_{\max}=\Phi(0)$ (Ref. 16).

Should particle natural oscillations thermalize between consecutive interactions with the field, the effect can be employed for particle cooling.[14]

The idea can be explained as follows. Assume, for simplicity, that Ω is constant. Then, since $\mathfrak{E}$ is conserved, the drift energy change for an initially free particle scattering off a ponderomotive barrier $[\Phi(\pm\infty)=0]$ equals

$$\Delta\mathcal{E}=\left(\frac{\omega}{\Omega}-1\right)\Delta\mathcal{W}; \tag{8}$$

hence, the sign of $\Delta\mathcal{E}$ is determined by the sign of $\Delta\mathcal{W}$ for given $\Delta\omega$. (For adiabatic interaction, both $\Delta\mathcal{E}$ and $\Delta\mathcal{W}$ would be zero.) Let us show that $\Delta\mathcal{W}>0$ for moderate $\eta\equiv T_0/\Phi_{\max}$, where T_0 is the initial temperature. At $\Lambda\equiv|1-\Omega/\omega|^{-1}\gg1$, one has $\Delta\mathcal{W}\sim\Lambda\Phi_{\max}$,[14,18] so $\Delta\mathcal{W}\gg T_0$, assuming $\Lambda\gg\eta$. Thus, $\Delta\mathcal{W}\approx\mathcal{W}(+\infty)>0$ for all particles, regardless of the initial value of $\mathcal{W}$. As a result, the condition $\omega<\Omega$ guarantees that $\Delta\mathcal{E}<0$, so all thermal particles are decelerated.

Suppose now that each particle encounters the field repeatedly, and the time between consecutive encounters exceeds the relaxation time of the particle natural oscillations. (Contrary to Refs. 1, the relaxation occurs outside the ac field here.) At each impact, the particle will lose about $|\omega/\Omega-1|\Delta\mathcal{W}\sim|\Phi|_{\max}$ of its drift energy $\mathcal{E}$, and yet get to the next encounter with negligible $\mathcal{W}$, as compared to $\Lambda|\Phi|_{\max}$. After about $\eta\gg1$ interactions, each particle will then be cooled down to $\mathcal{E}\sim|\Phi|_{\max}\ll T_0$.

Below this limit, the scattering automatically becomes adiabatic, meaning that further cooling is stopped, if $\hat{\epsilon}\ll1$, where $\hat{\epsilon}=\hat{\lambda}/L$, $\hat{\lambda}=\hat{v}/|\Delta\omega|$, and $\hat{v}\sim(|\Phi|_{\max}/m)^{1/2}$. At $\hat{\epsilon}\gtrsim1$, the cooling is impeded as well, though, because particle trapping may occur in this regime.[14] The trapping effect can be explained as follows. It is possible that, due to nonadiabatic deceleration, a particle can lose all of its kinetic energy $\mathcal{E}$ even before leaving the field. In this case, the particle will be bounced back by the decelerating slope of the barrier and may continue bouncing afterward; that is, become trapped (Fig. 3). The total number of bounce oscillations is limited due to phase space conservation. Yet if this number is large enough, the post-trapping dynamics of the particle correlates little with its pre-trapping dynamics, so the particle can randomly escape either forward or backward (Fig. 3). The overall scattering is then stochastic, and may lead to both "transmission" and "reflection." An analogy with quantum tunneling through a conservative field can be drawn in this case, as we also discussed in Sec. II.

VI. ONE-WAY WALLS AND CURRENT DRIVE

Resonant fields can form asymmetric barriers, or one-way walls, which transmit particles in one direction with a higher probability than in the other direction. The idea of a stationary asymmetric ponderomotive barrier was first proposed for magnetized charged particles[18,19,33] and can be explained as follows. In a mirror-type dc magnetic field $\mathbf{B}_0\approx\hat{\mathbf{z}}B_0(z)$,[34] the rf-produced ponderomotive potential Φ is inversely proportional to $\Delta\omega\equiv\omega-\Omega(z)$, where $\Omega=eB_0/mc$ is the gyrofrequency. As the cyclotron resonance is approached, $\Phi(z)$ grows unlimitedly, and the barrier remains repulsive at $\Delta\omega(z)>0$ and attractive at $\Delta\omega(z)<0$ if the rf field has a maximum at the location where $\Delta\omega(z)=0$ [Fig. 4(a)]. This property of the ponderomotive force also persists in the immediate vicinity of the resonance, where the adiabatic approximation fails.[18] As a result, the barrier accelerates and heats particles traveling in one direction, but adiabatically repels those traveling in the other direction, hence acting somewhat like a Maxwell demon. Contrary to the latter case, though, the minimum amount of cyclotron heating for transiting particles is nonzero here, as the barrier must conserve the particle phase space.[17] In Ref. 18, we propose how to adjust the field gradients to operate closer to this minimum.

Given an appropriate rf profile, a one-way wall effect can also be practiced in a uniform magnetic field. In Ref. 16, we show that the asymmetric field configuration depicted in Fig. 4(b) transmits most of thermal particles incident from the right ($L_2\ll\lambda$) but adiabatically reflects those incident from the left ($L_1\gg\lambda$) if $\Phi_{\max}\gtrsim T_0$ [Figs. 5(a) and 5(b)]. Since this technique does not require that Ω is a function of z, it can be more easily extrapolated to neutral particles such as atoms and molecules. Similar ideas (yet employing a

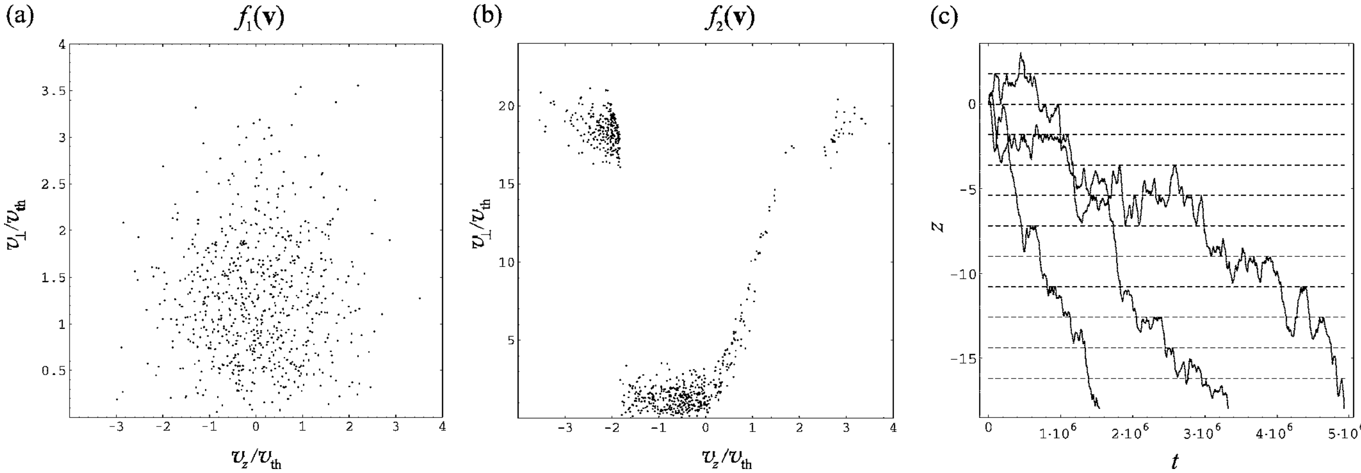

FIG. 5. (a) and (b) Velocity mapping produced by a ponderomotive barrier depicted in Fig. 4(b): (a) Gaussian distribution before scattering $f_1(\mathbf{v})$ and (b) distribution after scattering $f_2(\mathbf{v})$. Velocity is measured in units $v_{th}=\sqrt{T_0/m}$; $\Lambda=100$; $L_1=4\hat{\lambda}$; $L_2=0.02\hat{\lambda}$; $T=0.3\Phi_{max}$. For most of thermal particles, the mapping is deterministic, except for a small fraction of those nonadiabatically scattering off an abrupt slope of the barrier; about 80% of all particles traveling both ways end up on the left side of the barrier. (c) Sample trajectories of particles moving through a chain of such barriers (arbitrary units). Locations of individual barriers are shown by dashed lines; $\Delta L \gg L_1, L_2$.

bichromatic drive) have been proposed recently in the optical frequency range for manipulating atoms by laser radiation.[35,36] The applications envisioned for these techniques include atomic cooling via moving one-way walls.[35,37] Selective resonant properties of asymmetric barriers permit using those as well for isotope separation and related applications.

Unlike conservative barriers, ponderomotive one-way walls can also operate as emf sources and drive currents (or neutral flows) via introducing the anisotropy to otherwise thermal media.[19,33] For example, a chain of asymmetric barriers represents a new form of what is known as a ratchet potential.[16,38] Assuming that the barriers are separated by the distance ΔL large compared to the mean free path l_{mfp}, thermal particles will exhibit the average velocity $\langle v_d\rangle \sim (l_{mfp}/\Delta L)\sqrt{T_0/m}$ as they travel through the chain [Fig. 5(c)]; as a result, a current is generated.

In a toroidal system such as a tokamak, even a single one-way wall is sufficient to drive a current. In Ref. 19, we show that the efficiency of the asymmetric ponderomotive current drive (the current per power spent) can be comparable to or even larger than that of traditional schemes.[39] In Ref. 18, we also propose different field configurations to achieve the desirable operational states for the suggested technique.

VII. CONCLUSIONS

In summary, we offer new tools for manipulating charged and neutral particles by means of nonadiabatic ponderomotive forces and advance the analytical treatment of such forces. We show that in strongly inhomogeneous fields the ponderomotive force is phase dependent, and the particle dynamics resembles that of a quantum object in a conservative barrier. Probabilistic transmission through a ponderomotive potential is then possible and can be used for particle beam slicing. We also show that resonant fields can cool and trap particles exhibiting natural oscillations (e.g., Larmor rotation), as well as transmit them asymmetrically, hence acting as one-way walls. We find an approximate integral of particle motion for this case and introduce a new ponderomotive potential accordingly. Although only classical interactions are considered, the reported results, in principle, allow extrapolation to the quantum domain and can apply to such particles as atoms and molecules.

ACKNOWLEDGMENTS

This work was supported by DOE Contract No. DE-FG02-05ER54838, by the NNSA under the SSAA Program through DOE Research Grant No. DE-FG52-04NA00139, and by the National Science Foundation under Grant No. PHY99-07949.

[1] S. Chu, Rev. Mod. Phys. **70**, 685 (1998); C. N. Cohen-Tannoudji, Rev. Mod. Phys. **70**, 707 (1998); W. D. Phillips, Rev. Mod. Phys. **70**, 721 (1998).
[2] A. Ashkin, Proc. Natl. Acad. Sci. U.S.A. **94**, 4853 (1997); T. Iida and H. Ishihara, Phys. Rev. Lett. **90**, 057403 (2003); T. Iida and H. Ishihara, IEICE Trans. Electron. **E88-C**, 1809 (2005); T. Sugiura, T. Okada, Y. Inouye, O. Nakamura, and S. Kawata, Opt. Lett. **22**, 1663 (1997); S. Ito, H. Yoshikawa, and H. Masuhara, Appl. Phys. Lett. **78**, 2566 (2001), and references therein.
[3] A. Ashkin, Science **210**, 4474 (1980).
[4] In principle, the effects reported here also allow extrapolation to the quantum domain and can apply to such particles as atoms and molecules. See, e.g., Sec. VI.
[5] A. V. Gaponov and M. A. Miller, Sov. Phys. JETP **7**, 168 (1958).
[6] H. Motz and C. J. H. Watson, Adv. Electron. **23**, 153 (1967).
[7] R. R. Agayan, F. Gittes, R. Kopelman, and C. F. Schmidt, Appl. Opt. **41**, 2318 (2002), and references therein.
[8] Higher-order corrections apply to the expression for Φ [Eq. (1)] in stronger fields (Ref. 13).
[9] J. R. Cary and A. N. Kaufman, Phys. Rev. Lett. **39**, 402 (1977).
[10] T. Consoli and R. B. Hall, Nucl. Fusion **3**, 237 (1963); H. P. Eubank, Phys. Fluids **12**, 234 (1969); T. Hatori and T. Watanabe, Nucl. Fusion **15**, 143 (1975); A. J. Lichtenberg and H. L. Berk, Nucl. Fusion **15**, 999 (1975); G. Dimonte, B. M. Lamb, and G. J. Morales, Plasma Phys. **25**, 713 (1983); B. M. Lamb, G. Dimonte, and G. J. Morales, Phys. Fluids **27**, 1401 (1984).
[11] J. R. Ferron, N. Hershkowitz, R. A. Breun, S. N. Golovato, and R. Goulding, Phys. Rev. Lett. **51**, 1955 (1983); Y. Yasaka and R. Itatani, Phys.

Rev. Lett. **56**, 2811 (1986); P. L. Similon and A. N. Kaufman, Phys. Rev. Lett. **53**, 1061 (1984); P. L. Similon, Phys. Rev. Lett. **58**, 495 (1987).
[12] T. Watanabe, T. Watari, and N. Ohyabu, Fusion Technol. **27**, 515, Suppl. S (1995); S. Masuzaki, N. Ohno, and S. Takamura, J. Nucl. Mater. **220**, 1112 (1995); S. Masuzaki and S. Takamura, Contrib. Plasma Phys. **34**, 318 (1994); T. Watari, R. Kumazawa, T. Mutoh, T. Seki, K. Nishimura, and F. Shimpo, Nucl. Fusion **33**, 1635 (1993); A. Grossman, L. Schmitz, F. Najmabadi, and R. W. Conn, J. Nucl. Mater. **196**, 775 (1992).
[13] I. Y. Dodin and N. J. Fisch, Phys. Rev. E **74**, 056404 (2006).
[14] I. Y. Dodin and N. J. Fisch, Phys. Lett. A **349**, 356 (2006).
[15] I. Y. Dodin and N. J. Fisch, Phys. Rev. Lett. **95**, 115001 (2005).
[16] I. Y. Dodin and N. J. Fisch, Phys. Rev. E **72**, 046602 (2005).
[17] I. Y. Dodin and N. J. Fisch, Phys. Lett. A **341**, 187 (2005).
[18] I. Y. Dodin, N. J. Fisch, and J. M. Rax, Phys. Plasmas **11**, 5046 (2004).
[19] N. J. Fisch, J. M. Rax, and I. Y. Dodin, Phys. Rev. Lett. **91**, 205004 (2003); **93**, 059902(E) (2004).
[20] I. Y. Dodin and N. J. Fisch, J. Plasma Phys. **71**, 289 (2005).
[21] I. Y. Dodin, N. J. Fisch, and G. M. Fraiman, JETP Lett. **78**, 202 (2003).
[22] I. Y. Dodin and N. J. Fisch, Phys. Rev. E **68**, 056402 (2003).
[23] A. J. Lichtenberg and M. A. Lieberman, *Regular and Chaotic Dynamics*, 2nd ed. (Springer, New York, 1992), pp. 144–147.
[24] Particle motion in a dc magnetic field can also be treated similarly. See, e.g., Ref. 23, pp. 87–93, 99–100.
[25] R. K. Varma, Phys. Rep. **378**, 301 (2003).
[26] C. S. Unnikrishnan, Phys. Rev. E **70**, 028501 (2004).
[27] More precisely, near those corresponding to the elliptic points of the Poincaré mapping. See, e.g., Ref. 23, pp. 183–188.
[28] The particle oscillation period may not equal the field period (Sec. IV), so, in general, ϵ and ς are not the same.
[29] S. P. D. Mangles, C. D. Murphy, Z. Najmudin *et al.*, Nature **431**, 535 (2004); C. G. R. Geddes, C. Toth, J. van Tilborg, E. Esarey, C. B. Schroeder, D. Bruhwiler, C. Nieter, J. Cary, and W. P. Leemans, Nature **431**, 538 (2004); J. Faure, Y. Glinec, A. Pukhov, S. Kiselev, S. Gordienko, E. Lefebvre, J. P. Rousseau, F. Burgy, and V. Malka, Nature **431**, 541 (2004).
[30] M. D. Perry, D. Pennington, B. C. Stuart *et al.*, Opt. Lett. **24**, 160 (1999); M. H. Key, M. D. Cable, T. E. Cowan *et al.*, Phys. Plasmas **5**, 1966 (1998); S. P. Hatchett, C. G. Brown, T. E. Cowan *et al.*, Phys. Plasmas **7**, 2076 (2000); S.-W. Bahk, P. Rousseau, T. A. Planchon, V. Chvykov, G. Kalintchenko, A. Maksimchuk, G. A. Mourou, and V. Yanovsky, Opt. Lett. **29**, 2837 (2004).
[31] N. Naumova, I. Sokolov, J. Nees, A. Maksimchuk, V. Yanovsky, and G. Mourou, Phys. Rev. Lett. **93**, 195003 (2004); N. M. Naumova, J. A. Nees, I. V. Sokolov, B. Hou, and G. A. Mourou, Phys. Rev. Lett. **92**, 063902 (2004); J. Nees, N. Naumova, E. Power *et al.*, J. Mod. Opt. **52**, 305 (2005).
[32] B. Hafizi and R. E. Aamodt, Phys. Fluids **30**, 3059 (1987).
[33] E. V. Suvorov and M. D. Tokman, Sov. J. Plasma Phys. **14**, 557 (1988); A. G. Litvak, A. M. Sergeev, E. V. Suvorov, M. D. Tokman, and I. V. Khazanov, Phys. Fluids B **5**, 4347 (1993).
[34] $\boldsymbol{\nabla}\cdot\mathbf{B}=0$ is provided due to the transverse component $B_\perp \ll B$. For details, see, e.g., Eqs. (7) and (8) in Ref. 18.
[35] M. G. Raizen, A. M. Dudarev, Q. Niu, and N. J. Fisch, Phys. Rev. Lett. **94**, 053003 (2005).
[36] A. Ruschhaupt and J. G. Muga, Phys. Rev. A **70**, 061604(R) (2004).
[37] A. M. Dudarev, M. Marder, Q. Niu, N. J. Fisch, and M. G. Raizen, Europhys. Lett. **70**, 761 (2005).
[38] See, e.g., P. Reimann, Phys. Rep. **361**, 57 (2002); P. Hänggi, F. Marchesoni, and F. Nori, Ann. Phys. **14**, 51 (2005); F. Julicher, A. Ajdari, and J. Prost, Rev. Mod. Phys. **69**, 1269 (1997), and references therein.
[39] N. J. Fisch, Rev. Mod. Phys. **59**, 175 (1987).