# Princeton Plasma Physics Laboratory

PPPL-4038

PPPL-4038

# Quantized Eigenstates of a Classical Particle in a Ponderomotive Potential

I.Y. Dodin and N.J. Fisch

December 2004

PRINCETON PLASMA PHYSICS LABORATORY

Prepared for the U.S. Department of Energy under Contract DE-AC02-76CH03073.

# PPPL Report Disclaimers

## Full Legal Disclaimer



## Trademark Disclaimer



# PPPL Report Availability

This report is posted on the U.S. Department of Energy's Princeton Plasma Physics Laboratory Publications and Reports web site in Fiscal Year 2005. The home page for PPPL Reports and Publications is: http://www.pppl.gov/pub_report/

## Office of Scientific and Technical Information (OSTI):

Available electronically at: http://www.osti.gov/bridge.

Available for a processing fee to U.S. Department of Energy and its contractors, in paper from:

U.S. Department of Energy
Office of Scientific and Technical Information
P.O. Box 62
Oak Ridge, TN 37831-0062

Telephone: (865) 576-8401

Fax: (865) 576-5728

E-mail: reports@adonis.osti.gov

## National Technical Information Service (NTIS):

This report is available for sale to the general public from:

U.S. Department of Commerce
National Technical Information Service
5285 Port Royal Road
Springfield, VA 22161

Telephone: (800) 553-6847

Fax: (703) 605-6900

Email: orders@ntis.fedworld.gov

Online ordering: http://www.ntis.gov/ordering.htm

# Quantized eigenstates of a classical particle in a ponderomotive potential

I.Y. Dodin and N.J. Fisch
*Princeton Plasma Physics Laboratory, Princeton, NJ 08543*
(Dated: November 12, 2004)

The average dynamics of a classical particle under the action of a high-frequency radiation resembles quantum particle motion in a conservative field with an effective de Broglie wavelength $\lambda$ equal to the particle average displacement on a period of oscillations. In a "quasi-classical" field, with a spatial scale large compared to $\lambda$, the guiding center motion is adiabatic. Otherwise, a particle exhibits quantized eigenstates in a ponderomotive potential well, can tunnel through classically forbidden regions and experience reflection from an attractive potential. Discrete energy levels are also found for a "crystal" formed by multiple ponderomotive barriers.

PACS numbers: 52.35.Mw, 03.65.-w, 52.25.Dg

Under intense high-frequency (rf) radiation, a charged particle undergoes fast oscillations superimposed on the average drift motion. If the particle drift displacement $\lambda$ on a period of these oscillations is small, that is if

$$\lambda \ll L, \tag{1}$$

where $L$ defines the spatial scale of the field, the particle dynamics remains adiabatic. The average effect of the rf drive can then be approximated with an effective potential, known as the ponderomotive, or Miller potential [1]. This paper shows that at larger $\lambda$, when the approximation of a conservative force becomes invalid, the motion of a guiding center resembles dynamics of a *quantum* object, unlike the *classical* particle adiabatic motion. For a related problem, namely, the guiding center dynamics in a nonuniform magnetic field, this analogy was previously drawn by Varma, as reviewed in Ref. [2]. However, the explanation of the "macroquantum" effects in terms of a quantum-mechanical wave function remains controversial [3].

In our paper, in contrast to quantum approach [2], we show that purely classical systems exhibit quantum-*like* effects. Compared to the classic quantum problem, where the particle is a wave packet and the potential is sharply defined, here, the true particle is a point object, but the "potential" is a wave packet. However, we show that a guiding center can be treated as a quantum-like object, to which the sharply defined Miller force (rather than the true Lorentz force) is applied. In particular, we show that an rf-driven particle exhibits quantized eigenstates in a ponderomotive potential well, can tunnel through "classically forbidden" regions, and can experience reflection from an attractive potential. We also show that discrete energy levels exist in a "crystal" formed by multiple ponderomotive barriers.

For simplicity, consider an rf field applied over a uniform dc magnetic field, so that the adiabatic potential is given by [1, 4]

$$\Phi = \frac{e^2 E^2}{4m(\omega^2 - \Omega^2)}. \tag{2}$$

Here we assume a linearly polarized rf field $\mathbf{E}_{\rm rf} = \mathbf{x}^0 E(z) \sin\omega t$ perpendicular to the magnetic field $\mathbf{B}_0 = \mathbf{z}^0 B_0$, with gyrofrequency $\Omega = eB_0/mc$ close to the frequency of the rf drive, *i.e.*, $\Lambda \equiv (\omega/\Omega - 1)^{-2} \gg 1$. In a potential (2), the particle guiding center experiences conservative motion if

$$\lambda = 2\pi v_z/(\omega - \Omega) \tag{3}$$

remains small compared to the characteristic scale $L$ of the rf field profile $E(z)$, as required by (1). In this case, averaged over the rf period $2\pi/\omega$, the Larmor period $2\pi/\Omega$, and the beat period $2\pi/(\omega-\Omega)$, two adiabatic invariants of particle motion are conserved. Those are the magnetic moment of free Larmor rotation $\mu = m(\mathbf{v}_\perp - \mathbf{v}_\sim)^2/2B_0$ and the quasi-energy of the particle motion along a dc magnetic field line $\mathfrak{E} = \frac{1}{2}mv_z^2 + \mu B_0 + \Phi$. (Here $\mathbf{v}$ is the particle velocity, and $\mathbf{v}_\sim$ is the velocity of rf-induced particle oscillations transverse to $\mathbf{B}_0$.)

Suppose now a particle confined by the potential (2) in the vicinity of its local minimum. In the adiabatic approximation, such a particle experiences conservative bounce motion described by the equation $\ddot{z} = -\Phi'(z)$. However, bounce oscillations remain conservative only approximately for finite $\lambda/L$. Each time a particle bounces off a ponderomotive wall, it either gains or loses energy, depending on its velocity $\mathbf{v}_0$ at the bottom of the well. (Similar nonadiabatic energy exchange was studied in a number of theoretical [5, 6] and experimental [7, 8] papers.) If for some $\mathbf{v}_0$ the energy gain over the bounce period is precisely zero no matter what the phase of the particle, the particle can be thought of as occupying a stationary *eigenstate* of a ponderomotive well, as its motion remains strictly periodic regardless of nonadiabaticity of interaction with the rf field.

In this paper, we report analytic and numerical identification of such eigenstates and describe their structure in the simplest case of an even potential $\Phi(z)$ in a subclass of bounce oscillations with $v_\perp(z=0) = 0$. To do so, let us first introduce the following dimensionless notation. We measure the frequency $\omega$ in units $\Omega$, the time – in units $\Omega^{-1}$, the particle velocity – in units $c$, the spatial coordinates – in units $c/\Omega$, and the electric field – in units $mc\Omega/e$. Solving for the perpendicular motion

$$\dot{w} + iw = E\,\sin\omega t, \tag{4}$$

where $w = v_x + iv_y$, one gets $w = \psi e^{-it}$, where, to the leading order in $\Lambda$,

$$\psi = \frac{i}{2}\int_0^z E(\tilde{z})\, e^{-i\chi(\tilde{z})}\, \frac{d\tilde{z}}{v_z(\tilde{z})}, \qquad (5)$$

and $\chi(z) = (\omega - 1)\int_0^z d\tilde{z}/v_z(\tilde{z})$. Averaging over the frequency $\omega \approx \Omega$, we get [5] for the longitudinal motion:

$$\ddot{z} = -\Phi'(z) - \phi'(z), \qquad (6)$$

where $\phi(z)$ is a quasi-potential vanishing in the adiabatic limit (1):

$$\phi(z) = -\frac{1}{8(\omega-1)}\left|\int_0^z E'(\tilde{z})\, e^{-i\chi(\tilde{z})}\, d\tilde{z}\right|^2. \qquad (7)$$

For clarity, assume $E(z=0) = 0$ and $\omega > 1$ ($\omega > \Omega$ in dimensional units). A particle starting with $v_z = v_0$ at the bottom of the well $z = 0$ will be decelerated by the rf field, come to a stop at a turning point $z(t = t_\mathrm{s}) = A$, and be reflected backward. Since particle motion may be nonadiabatic, the trajectory after reflection generally will not be symmetric to that before the reflection. However, it might be possible that the symmetry does exist for some $v_0$. Having such a case requires that $\phi(z(t))$ is an even function of $t - t_\mathrm{s}$, which is equivalent to

$$\arg\psi(A) + \chi(A) = \pi n, \qquad (8)$$

where $n$ is an integer. If this condition is satisfied, the particle is returned to $z = 0$ with the longitudinal and transverse energies precisely matching their initial values $\mathcal{E}_{||,0} = \frac{1}{2}v_0^2$, $\mathcal{E}_{\perp,0} = 0$. If $\Phi(z)$ is even, the phase-space trajectory of such a particle will form a *closed loop* on the plane $(z, v_z)$ (Fig. 1). In contrast to other particles, those with $v_0$ satisfying the condition (8) are not heated by the rf field despite undergoing nonadiabatic motion.

Eigenstates of a ponderomotive well are quite similar to those of a quantum particle in a true potential. Indeed, a particle with a proper $v_0$ will remain on a stationary "energy level" regardless of the initial rf phase. Hence, for an ensemble of particles with equal $v_0$ but different phases, the structure of an eigenstate will determine the probability of finding a particle at a certain location $z$. Extending the analogy, let us simplify the quantization rule (8) to the "quasi-classical" limit. At $n \gg 1$ Eq. (8) yields an approximate solution $(\omega - 1)t_\mathrm{s} \approx \pi n$. For a particle at a stationary eigenstate, $v_z$ can be considered a single-valued function of $z$. Hence, extending the integration from a quarter to the full bouncing period, the quasi-classical limit of (8) reads

$$\oint k\, dz = 4\pi n, \qquad (9)$$

where $k = 2\pi/\lambda$. Eq. (9) is analogous to the Bohr-Sommerfeld quantization condition for even energy levels $\mathcal{E}_{2n}$ of a quantum particle in a potential well. (Also, analogously, at $n \gg 1$ the velocity $v_z$ can be approximately calculated according to the adiabatic model.) Thus, to allow stationary bounce motion in an rf field, the amplitude of particle oscillations $A$ may not be less than the effective de Broglie wavelength $\lambda_0 = \lambda(v_0)$. Assuming that the field has a spatial scale $L$, the total number of levels can be estimated as $L/\lambda_0$, and the ground energy level $\mathcal{E}_1$ can be expected at $A \sim \lambda_0$, hence satisfying the order-of-magnitude equation $\mathcal{E}_1 \sim \Phi(\sqrt{\mathcal{E}_1\Lambda})$.

As an example, let us consider $E(z) = q|z|^\alpha$, $q = \mathrm{const}$, $\alpha > 0$. By approximating $z(t)$ with a parabolic function on a half of a bounce period, one can rewrite Eq. (9) in a more precise form $\oint k\, dz = 4\pi\,(n + \frac{1}{2}\,\alpha)$. Correspondingly, the "quasi-classical" energy spectrum is given by

$$\mathcal{E}_n = \hat{\mathcal{E}}\,\left(n + \frac{\alpha}{2}\right)^{2\alpha/(1-\alpha)}, \qquad (10)$$

where $\hat{\mathcal{E}}$ depends solely on the parameters of the field:

$$\hat{\mathcal{E}} = \frac{1}{2}\left\{\frac{q\Lambda^{\frac{1}{4}}}{2}\left[\sqrt{\pi\Lambda}\,\frac{\Gamma(\frac{1}{2}+\frac{1}{2\alpha})}{\Gamma(1+\frac{1}{2\alpha})}\right]^\alpha\right\}^{2/(1-\alpha)}. \qquad (11)$$

Hence, $\mathcal{E}_n$ increases with $n$ if $\alpha < 1$. In this case, all the energy levels lie above $\hat{\mathcal{E}}$, and the quasi-classical limit is approached as $v_0 \to \infty$. On the contrary, when $\alpha > 1$, all the eigenstates have energies below $\hat{\mathcal{E}}$ and become quasi-classical as $v_0 \to 0$. (At $\alpha = 1$, which would correspond to a linear pendulum in the adiabatic limit, degeneracy is observed: in this case all trajectories are self-similar, and the time $t_\mathrm{s}$ is independent of $v_0$.) Such a difference between the two cases is due to the fact that a profile $E(z) \propto |z|^\alpha$ does not have an intrinsic spatial scale, which is thus effectively determined by the amplitude of the particle bounce oscillations $A \propto v_0^{1/\alpha}$. The adiabaticity condition (1) with $\lambda$ given by (3) then can be put as $(v_0/\hat{v})^{\alpha-1} \ll 1$ (where $\hat{v} = \sqrt{\hat{\mathcal{E}}}$), which is satisfied for $v_0 \gg \hat{v}$ if $\alpha < 1$ and $v_0 \ll \hat{v}$ if $\alpha > 1$ (Fig. 2).

Quantization of stationary energy levels of particle motion confined by a ponderomotive force resembles that of a quantum particle in a potential well and has a similar nature. Analogously to a quantum object, the particle guiding center is not a zero-dimensional entity but can be assigned a *phase*, which is the phase of a real particle oscillations in an rf field. The distance $\lambda$, which determines the "uncertainty" of the guiding center location, can be naturally treated as the effective de Broglie wavelength of the latter, and the approximation of a local Miller force remains applicable only if the guiding center coordinates are well defined. What is remarkable is that this "quantum" analogy can be extended further and applies also to a freely moving (non-confined) particle. Indeed, the longitudinal force acting on a guiding center is the average Lorentz force proportional to the transverse particle velocity. If a particle incident on a localized rf barrier is fast enough, as given by (1), it will not have sufficient time to gain transverse energy from the field. Hence, it will neither experience significant ponderomotive acceleration. Such a particle will then be able to

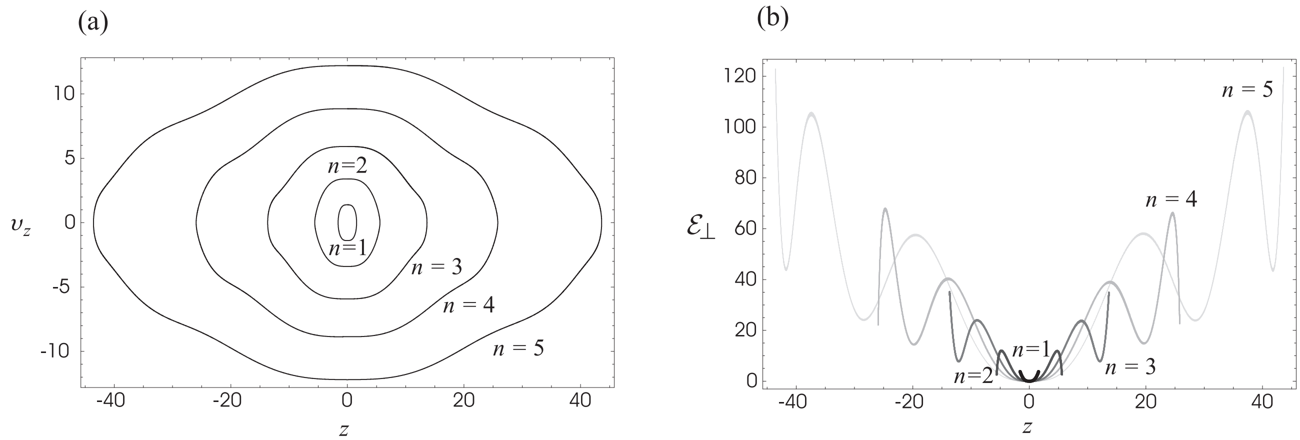


FIG. 1: First five stationary eigenstates of a guiding center trapped within a ponderomotive potential formed by an rf field with the amplitude $E(z) = q|z|^\alpha$ with $\alpha = 0.6$, $q = 10^{-3}$: (a) phase plane $(z, v_z)$, (b) perpendicular energy $\mathcal{E}_\perp(z)$ ($v_z$ is measured in units $\hat{v} = \sqrt{\hat{\mathcal{E}}}$; $\mathcal{E}_\perp(z)$ is measured in units $\hat{\mathcal{E}}\sqrt{\Lambda}$ [see Eq. (11)]).

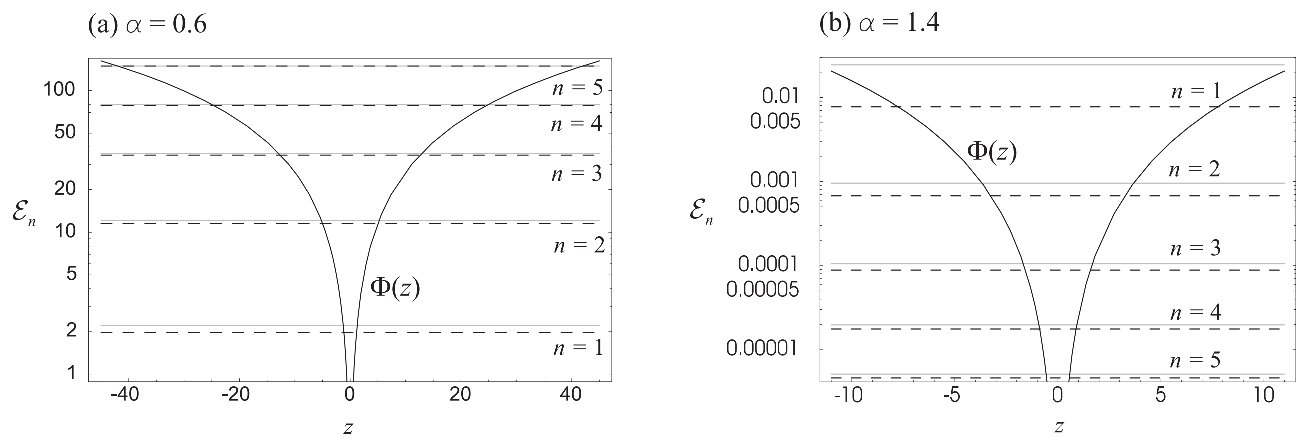


FIG. 2: First five energy levels $\mathcal{E}_n$ [dashed – numerical; solid gray – quasi-classical, given by Eq. (10)] of a guiding center bouncing within a ponderomotive potential $\Phi(z)$ (solid black) formed by an rf field with the amplitude $E(z) = q|z|^\alpha$: (a) $\alpha = 0.6$, $q = 10^{-3}$; (b) $\alpha = 1.4$, $q = 2 \times 10^{-4}$ ($\mathcal{E}_n$ and $\Phi$ are measured in units $\hat{\mathcal{E}}$).

penetrate ("tunnel") through the "classically forbidden" region $\frac{1}{2}\,v_0^2 < \Phi(z)$, just like a quantum particle having a de Broglie wavelength of the order of the field scale.

The depression of a ponderomotive force (in comparison with the adiabatic model) in the vicinity of a cyclotron resonance is confirmed in our numerical calculations, and was reported also in Refs. [7]. What has not been reported yet is that "quantum" properties are also inherent to attractive ponderomotive barriers ($\Phi < 0$), which appear to be capable of reflecting particles. First, note that Eq. (6) yields a theorem [5]

$$\Delta\mathcal{E}_{||} = (\omega - 1)\Delta\mathcal{E}_\perp, \tag{12}$$

which connects the integral changes of longitudinal and transverse energies of a particle at $t \to \infty$. In the adiabatic limit, the energy change (12) remains exponentially small with respect to $\epsilon \equiv \lambda_0/2\pi L$. Suppose though that $\epsilon \gtrsim 1$ and $\Phi < 0$, meaning that $\omega < 1$. Since $\Delta\mathcal{E}_\perp > 0$, in this case a particle *loses* $\mathcal{E}_{||}$ as a result of interaction. At some $v_0 = v_z(t \to -\infty)$, this deceleration can become sufficient to *trap* a particle in a potential well: a particle entering the rf field freely can be decelerated inside and bounced back toward the stronger field at the exit (Fig. 3). It can be shown that the trapping condition (also checked numerically) can be written as $\hat{\epsilon} \equiv \hat{\lambda}/2\pi L \gtrsim 1$. That is, if $\hat{\epsilon} \gtrsim 1$, at sufficiently small $v_0$ a particle may be trapped within a potential well, but if $\hat{\epsilon} \lesssim 1$, trapping is impossible regardless of $v_0$. This can be explained as follows. Slow particles ($v_0 \ll \hat{v}$) are accelerated ponderomotively up to the velocity of the order of $\hat{v}$. If the latter itself is large enough ($\hat{\epsilon} \gtrsim 1$), nonadiabatic effects have to reveal for all, even initially slow particles, some of which may then experience trapping. On the contrary, at $\hat{\epsilon} \lesssim 1$, slow particles remain adiabatic and hence cannot be trapped. As for fast particles ($v_0 \gg \hat{v}$), in both cases they have enough energy to overcome the ponderomotive deceleration and avoid trapping.

What is the "destiny" of a "once trapped" particle?

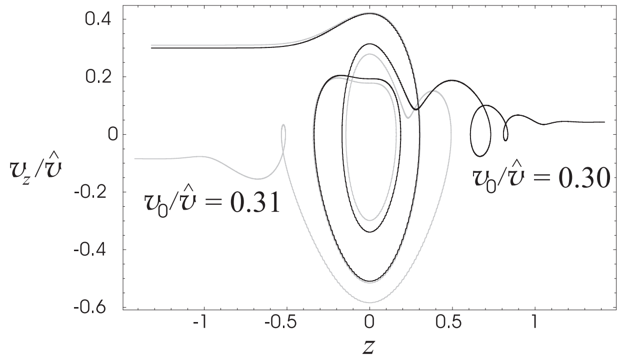


FIG. 3: Longitudinal velocity $v_z$ vs $z$ for a particle being trapped and released by an attractive ponderomotive potential with $\hat{\epsilon} = 3$ ($v_z$ is measured in units $\hat{v}$; $E_0 = 0.001$, $\Lambda = 10^4$, $L = 0.33$; the initial magnetic moment is equal to zero): $v_z = 0.30\,\hat{v}$ (black) and $v_z = 0.31\,\hat{v}$ (gray).

Because of the phase space conservation requirement, particles may not stay trapped forever. However, if the number of bounce oscillations within a potential well is large, the post-trapping dynamics of a particle correlates little with its pre-trapping dynamics. Hence the direction, to which the particle is released, is almost uncorrelated with the initial velocity (Fig. 3), and a particle can escape toward the direction opposite to $v_0$, which qualifies as *reflection*. The effect disappears under the condition (1), and again resembles a quantum phenomenon when a particle can be reflected by an attractive potential if having de Broglie wavelength of the order of $L$.

Because of clearly stochastic behavior inside a ponderomotive well, a particle traveling through a chain of such potentials would undergo a random walk, as each of the potentials can scatter a particle back and forth with roughly equal probability. Hence, a sufficiently long chain of ponderomotive barriers violating the condition (1) acts like a diffusive mirror. However, among $v_0$, for which stochastic dynamics is realized, there exists a countable set of regular trajectories, at which a particle can "collisionlessly" travel through a "crystal" formed by multiple barriers. These trajectories can then be attributed as stationary eigenstates of a free particle moving in a "ponderomotive crystal", which has a well-known quantum analogue in the physics of solid state. Such eigenstates (Fig. 4) we found numerically for chains of both attractive and repulsive potentials. Similarly to bounce oscillations within a potential well, the ground energy level ($n = 1$) of a transmitting particle is located at $v_0 \sim \hat{v} \sim \sqrt{|\Phi|_{\max}}$. Higher levels ($n > 1$) are located at larger energies, and at $n \gg 1$ (corresponding to $v_0 \gg \hat{v}$) the particle motion becomes "classical", that is, in this case, only slightly disturbed by the ponderomotive force.

The quantum-like effects run counter to what follows from the adiabatic theory predicting regular deterministic dynamics to all orders of $\epsilon$. This is due to the fact that such elaboration is done via asymptotic and, hence, approximate methods, such as Lie techniques [9]. While the limited nature of these methods is often failed to mention, it is crucial for understanding particle dynamics at finite $\epsilon$. Therefore, in addition to the academic interest in the showing that very general classical systems exhibit quantum effects, capturing the effects we describe here will be a challenge to the existing computational and analytical techniques in plasma kinetic theory.

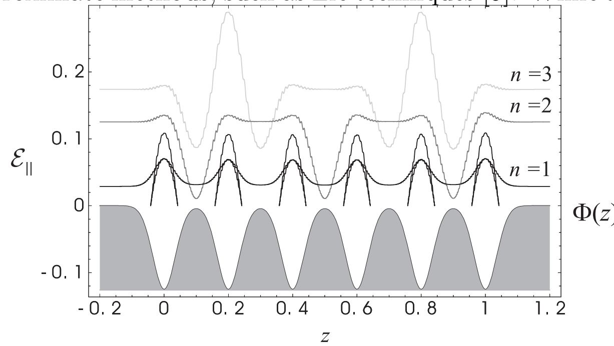


FIG. 4: First three stationary eigenstates of a free particle traveling through a "crystal" formed of multiple ponderomotive barriers of Gaussian shape $E = E_0\,\exp(-z^2/L^2)$. Shown is $\Phi(z)$ (shaded) and the longitudinal energy $\mathcal{E}_{||}$ vs $z$ (solid) (energy units are $\hat{v}^2$; $E_0 = 0.001$, $\Lambda = 10^3$, $L = 0.05$, the distance between the individual barriers is $8L$).

The work is supported by DOE contract DE-AC0276-CHO3073.

[1] A. V. Gaponov and M. A. Miller, Sov. Phys. JETP **7**, 168 (1958); H. Motz and C. J. H. Watson, Advances in Electronics **23**, 153 (1967).
[2] R. K. Varma, Phys. Reports **378**, 301 (2003).
[3] C. S. Unnikrishnan, Phys. Rev. E **70**, 028501 (2004).
[4] I. Y. Dodin and N. J. Fisch, to appear in J. Plasma Phys. (2004).
[5] I. Y. Dodin, N. J. Fisch, and J. M. Rax, Phys. Plasmas **11**, 5046 (2004).
[6] N. J. Fisch, J. M. Rax, and I. Y. Dodin, Phys. Rev. Lett. **91**, 205004 (2003); E. Canobbio, Nuclear Fusion **9**, 27 (1969); A. J. Lichtenberg and H. L. Berk, Nuclear Fusion **15**, 999 (1975); T. Consoli and R. B. Hall, Nuclear fusion **3**, 237 (1963); T. Hatori and T. Watanabe, Nuclear fusion **15**, 143 (1975).
[7] G. Dimonte, B. M. Lamb, and G. J. Morales, Plas. Phys. **25**, 713 (1983); B. M. Lamb, G. Dimonte, and G. J. Morales, Phys. Fluids **27**, 1401 (1984).
[8] H. P. Eubank, Phys. Fluids **12**, 234 (1969).
[9] A. J. Lichtenberg and M. A. Lieberman, *Regular and Chaotic Dynamics* (Springler Verlag, New York, 1992), 2nd ed.

## External Distribution

Plasma Research Laboratory, Australian National University, Australia
Professor I.R. Jones, Flinders University, Australia
Professor João Canalle, Instituto de Fisica DEQ/IF - UERJ, Brazil
Mr. Gerson O. Ludwig, Instituto Nacional de Pesquisas, Brazil
Dr. P.H. Sakanaka, Instituto Fisica, Brazil
The Librarian, Culham Laboratory, England
Mrs. S.A. Hutchinson, JET Library, England
Professor M.N. Bussac, Ecole Polytechnique, France
Librarian, Max-Planck-Institut für Plasmaphysik, Germany
Jolan Moldvai, Reports Library, Hungarian Academy of Sciences, Central Research Institute for Physics, Hungary
Dr. P. Kaw, Institute for Plasma Research, India
Ms. P.J. Pathak, Librarian, Institute for Plasma Research, India
Ms. Clelia De Palo, Associazione EURATOM-ENEA, Italy
Dr. G. Grosso, Instituto di Fisica del Plasma, Italy
Librarian, Naka Fusion Research Establishment, JAERI, Japan
Library, Laboratory for Complex Energy Processes, Institute for Advanced Study, Kyoto University, Japan
Research Information Center, National Institute for Fusion Science, Japan
Dr. O. Mitarai, Kyushu Tokai University, Japan
Dr. Jiangang Li, Institute of Plasma Physics, Chinese Academy of Sciences, People's Republic of China
Professor Yuping Huo, School of Physical Science and Technology, People's Republic of China
Library, Academia Sinica, Institute of Plasma Physics, People's Republic of China
Librarian, Institute of Physics, Chinese Academy of Sciences, People's Republic of China
Dr. S. Mirnov, TRINITI, Troitsk, Russian Federation, Russia
Dr. V.S. Strelkov, Kurchatov Institute, Russian Federation, Russia
Professor Peter Lukac, Katedra Fyziky Plazmy MFF UK, Mlynska dolina F-2, Komenskeho Univerzita, SK-842 15 Bratislava, Slovakia
Dr. G.S. Lee, Korea Basic Science Institute, South Korea
Institute for Plasma Research, University of Maryland, USA
Librarian, Fusion Energy Division, Oak Ridge National Laboratory, USA
Librarian, Institute of Fusion Studies, University of Texas, USA
Librarian, Magnetic Fusion Program, Lawrence Livermore National Laboratory, USA
Library, General Atomics, USA
Plasma Physics Group, Fusion Energy Research Program, University of California at San Diego, USA
Plasma Physics Library, Columbia University, USA
Alkesh Punjabi, Center for Fusion Research and Training, Hampton University, USA
Dr. W.M. Stacey, Fusion Research Center, Georgia Institute of Technology, USA
Dr. John Willis, U.S. Department of Energy, Office of Fusion Energy Sciences, USA
Mr. Paul H. Wright, Indianapolis, Indiana, USA

The Princeton Plasma Physics Laboratory is operated
by Princeton University under contract
with the U.S. Department of Energy.

Information Services
Princeton Plasma Physics Laboratory
P.O. Box 451
Princeton, NJ 08543

Phone: 609-243-2750
Fax: 609-243-2751
e-mail: pppl_info@pppl.gov
Internet Address: http://www.pppl.gov