Eur. Phys. J A **2**, 115–121 (1998)

THE EUROPEAN PHYSICAL JOURNAL A

Classically dynamical behaviour of a nucleon in heavy nuclei

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Reiceved: 10 November 1997 / Revised version: 27 January 1998
Communicated by W. Weise

Abstract. Within the framework of the two-center shell model the classically dynamical behaviour of a nucleon in heavy nuclei is investigated systematically with the change of nuclear shape parameters for the first time. It is found that as long as the nucleonic energy is appreciably higher than the height of the potential barrier there is a good quantum-classical correspondence of nucleonic regular (chaotic) motion. Thus, Bohigas, Giannoni and Schmit conjecture is confirmed once again. We find that the difference between the potential barrier for prolate nuclei and that for oblate ones is reponsible for the energy-dependence difference between the nucleonic chaotic dynamics for prolate nuclei and that for oblate ones. In addition, it is suggested that nuclear dissipation is shape-dependent, and strong nuclear dissipation can be expected for medium or large separations in the presence of a considerable neck deformation built on a pronounced octupole-like deformation, which provides us a dynamical understanding of nuclear shape dependence of nuclear dissipation.

PACS. 05.45.+b Theory and models of chaotic systems – 21.60.Cs Shell model – 03.20.+i Classical mechanics of discreta systems: general mathematical aspects – 11.10.Lm Nonlinear or nonlocal theories and models

1 Introduction

The classical nonlinear dynamics of Hamiltonian systems has been widely and deeply investigated, and the field of interest is being extended to quantum systems as well [1]. But since the Schrödinger equation governing the quantum mechanical motion is linear and time reversible, one can hardly connect it with nonintegrability [2]. In addition, the conventional method of pursuing a motion of a representative point in phase space by analyzing the trajectory is not possible for quantum system. Furthermore, for highly unstable quantum systems perturbation theory is no longer adequate. All of these reasons make it difficult to have a clear definition for quantum chaos. Nevertheless, Bohigas [3] found that during the transition of a classical analog of a quantum system from regular to chaotic motion, the quantum behaviour does not manifest itself in a specific energy level or quantum state, but in statistical fluctuation properties of the global energy levels of the system. The systems whose classical analogs are integrable show Piosson fluctuations, whereas the systems whose classical analogs are fully chaotic show fluctuation patterns of a Gaussian orthogonal ensemble (GOE). Since both the Poisson and GOE distribution functions contain no free parameters, in 1984, Bohigas, Giannoni and Schmit (BGS) conjectured that this phenomenon was generic [4]. From then on Seligman et al. [5] studied a two-dimensional Hamiltonian which classically shows a transition from regular to chaotic behaviour. Delande and Gay [6] studied numerically the hydrogen atom in a magnetic field, another system which shows a transition from regularity to chaos.

Wintgen and Marxer [7] examined the anisotropic Kepler problem, while Meredith et al. [8] and Xu et al. [9] considered the three-orbital Lipkin-Meshkov-Glick model. In each these cases, the analog of a classically regular system showed Poisson fluctuation and the analog of a classically chaotic system displayed GOE behaviour. Now it has been generally accepted that quantum systems are regular (chaotic) if their spectral fluctuations can be described by a Poisson (GOE) statistics.

The nearest neighbour spacing distribution P(s) and spectral rigidity $\Delta_3(L)$ are two commonly used statistical measures [1]. The former is equal to the probability density for two neighbouring levels having the spacing s and measures the degree of level repulsion. The latter is the least square deviation of the number of spacings in a given energy interval L from the best fit to a straight line, which signifies the long-range correlations of a quantum spectrum. For the GOE statistics, the spacing distribution

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$$P(s) = \frac{\pi}{2}\ s\ \exp(-\frac{\pi\ s^2}{4}), \quad (1)$$

and the $\Delta_3(L)$ can only be evaluated numerically, but it approaches the value

$$\Delta_3(L) \cong \frac{1}{\pi^2}\ (lnL - 0.0687), \quad (2)$$

for large L.

For the Poissonian statistics,

$$P(s) = \exp(-s), \quad (3)$$

$$\Delta_3(L) = \frac{L}{15}. \quad (4)$$

One of the best systems for the study of quantum chaos is the atomic nucleus. Roughly speaking, the studies of quantum chaos related to nuclei have two aspects. One is to discuss the chaos in realistic nuclei [1,10]. The other is to investigate the single-particle (nucleonic) motion in a deformed nucleus using mean field approach [11,13]. The symmetry (or absence of symmetry) of the mean field related to the geometric shape of nuclei determines the regularity (or chaoticity) of the single-particle motion. Usually, quadrupole deformation is considered to be the major deviation from spherical symmetry. However, more recently a possible octupole contribution has been taken into account for a number of reasons [14,15]. Inclusion of an octupole deformation in addition to a quadrupole deformation (neglecting the spin-orbit coupling and l^2-correction terms) was discussed by Heiss et al. [11]. They found that when an octupole deformation is switched on, the nucleonic motion is regular for prolate nuclei, however, chaotic for oblate ones. It was also shown that there is a good quantum-classical correspondence of nucleonic regular (chaotic) motion. The multipole deformation potential mentioned above corresponds to one-center shell model which does not contain the proper asympototic shapes, and is not suitable to describe the heavy nuclei with necked-in or very elongated shapes. Due to the inclusion of a neck deformation, the two-center shell model (TCSM) [16] can describe not only the small deformation but also the large deformation (including necked-in or very elongated ones).

In [17], in the framework of the TCSM we calculated the P(s) and $\Delta_3(L)$ of single-particle energy levels of heavy nuclei when the shape parameters of a nucleus are changed systematically. The shape parameter region in which both the P(s) and $\Delta_3(L)$ are approximately close to those of GOE has been found. We call the shape parameter region as the chaotic region of shape parameter, in which for a nucleon in heavy nuclei the quantum chaotic motion is realized. Although the relationship between the quantum chaotic (regular) motion and shape deformations has been clarified in [17], the classical dynamics of a nucleon in heavy nuclei has not been studied. The present paper is to provide, for the first time, a detailed analysis of the classically dynamical behaviour of a nucleon in heavy nuclei by means of the TCSM. Since the neck deformation being a specific feature of the TCSM plays an important role in heavy ion collision, fusion-fission and nuclear molecular states, its influence on the classical dynamics will be particularly paid attention. Our aim is to seek the connection between the classical dynamics and quantum spectral statistics of a nucleon in heavy nuclei to see whether the BGS conjecture is still valid for such a complex system. In addition, the possible relationship between the shape deformations and nuclear dissipation will be discussed.

2 The two-center shell model and the chaotic region of shape parameter

Neglecting the spin-orbit coupling and l^2-correction terms, the single-particle Hamiltonian of the TCSM in cylinder coordinetes z, ρ, ϕ is as follows

$$H = -\frac{\hbar^2\ \nabla^2}{2m_0} + V(\rho, z). \quad (5)$$

The potential reads

$$V(\rho,z) = \begin{cases} \frac{1}{2}m_0\omega_{z_1}^2 z'^2 + \frac{1}{2}m_0\omega_{\rho_1}^2\rho^2, & z < z_1 \\ \begin{array}{l}\frac{f_0}{2}m_0\omega_{z_1}^2\ z'^2\ (1 + c_1 z' + d_1 z'^2) \\ \quad + \frac{1}{2}m_0\omega_{\rho_1}^2\ (1 + g_1 z'^2)\rho^2\end{array}, & z_1 < z < 0 \\ \begin{array}{l}\frac{f_0}{2}m_0\omega_{z_2}^2\ z'^2\ (1 + c_2 z' + d_2 z'^2) \\ \quad + \frac{1}{2}m_0\omega_{\rho_2}^2 (1 + g_2 z'^2)\rho^2\end{array}, & 0 < z < z_2 \\ \frac{1}{2}m_0\omega_{z_2}^2 z'^2 + \frac{1}{2}m_0\omega_{\rho_2}^2\rho^2, & z > z_2 \end{cases} \quad (6)$$

by denoting the positions of the centers of the two fragments by z_1 and z_2, $z_1 \leq 0 \leq z_2$, with the abbreviation

$$z' = \begin{cases} z - z_1, & z < 0 \\ z - z_2. & z > 0 \end{cases}$$

All shape deformation parameters of a nucleus can be reduced to the following five independent ones. (a) The separation of the two centers $\Delta z = z_2 - z_1$. (b) The neck deformation of a nucleus $\epsilon = \frac{E_0}{E'}$, where $E' = \frac{1}{2}m_0\omega_{z_i}^2 z_i^2$ (i=1,2), E_0 is the actual height of the barrier [16]. ϵ=0 corresponds to ovaloids, ϵ=1 to well necked-in shapes. (c) The mass asymmetry $X_i = \frac{(A_1 - A_2)}{(A_1 + A_2)}$ which ranges from 0 to 1. A_1 and A_2 are the mass numbers of the fragments, which have no explicit expression and are evaluated numerically [16]. (d) The ellipsoidal deformations of the fragments (local deformations) $\beta_i = \frac{\omega_{\rho_i}}{\omega_{z_i}}$ (i=1,2). If X_i =0, β_1=1=β_2 and ϵ=0, the shape deformation is a pure quadrupole one. Provided $X_i \neq 0$, an octupole-like deformation appears. If X_i is larger, and local deformations β_1 and β_2 are quite asymmetry (i. e. one of the fragments appears pronounced oblate, the other the pronounced prolate.), then a larger octupole-like deformation is expected.

The parameters in (6) can be written in the form

$$f_0 = 4\epsilon, \quad c_i = \frac{1}{z_i}, \quad d_i = \frac{1}{4z_i^2} \quad (i = 1, 2),$$

$$g_1 = \frac{1-Q^2}{z_1 \Delta\ z},\ g_2 = \frac{1-Q^2}{Q^2 z_2 \Delta\ z},$$
$$z_1 = -\frac{Q_3 \Delta\ z}{1+Q_3},\ z_2 = \frac{\Delta\ z}{1+Q_3}.$$

with $Q = \frac{\omega_{\rho 2}}{\omega_{\rho_1}}$, $Q_3 = \frac{\beta_2}{\beta_1 Q}$. The frequencies ω_{z_i} and ω_{ρ_i} (i=1,2) can be determined by the five shape parameters together with the requirements of the volume conservation and smooth joining of the potential. Thus when the five shape parameters are given all of the parameters in (6) can be determined uniquely.

In [17], it is shown that at $X_i = 0, \beta_1 = 1.0 = \beta_2$ (equivalent to no octupole-like deformation) the P(s) and $\Delta_3(L)$ always appear to be the Poisson-types for all possible values of ϵ and Δz. While $X_i \neq 0$ and β_1, β_2 deviate from 1.0 (i. e. an octupole-like deformation is added), the P(s) and $\Delta_3(L)$ depart from the Poisson-types. It was found that for heavy nuclei the single-particle spacing distribution and spectral rigidity are approximately close to those of GOE when the shape parameters fall into the following region [17]:

$$2.0\ <\ \Delta\ z\ <\ 5.0\ fm,\ \ 0.4\ <\ X_i\ <\ 0.7,$$
$$0.2\ <\ \beta_1\ <\ 0.5,\ \ 2.0\ <\ \beta_2\ <\ 4.0,$$
$$\epsilon_{min}\ <\ \epsilon\ <\ 0.9,$$
$$(\epsilon_{min} = 0\ \ for\ \ 2.0\ <\ \Delta\ z < 3.0,$$
$$\epsilon_{min} = 0\ \sim\ 0.3\ \ for\ \ 3.0\ <\ \Delta\ z <\ 5.0\ fm).\quad (I)$$
$$5.0\ <\ \Delta\ z\ <\ 11.0\ fm,\quad 0.3\ <\ X_i\ <\ 0.8,$$
$$0.3\ <\ \beta_1\ <\ 0.6,\ \ 2.0\ <\ \beta_2\ <\ 5.0,$$
$$0.3\ <\ \epsilon\ <\ 0.8.\quad (II)$$

Both of the parts (I) and (II) of the chaotic region exhibit that the appearance of the quantum chaotic motion requires a pronounced octupole-like deformation. The shape parameters in (I) and (II) correspond to the oblate and prolate shapes of a nucleus, respectively. Therefore the quantum chaotic motion occurs not only in the heavy nuclei with oblate shapes but also in prolate ones with a considerable neck deformation. The neck deformation appreciably influences the occurence of the quantum chaotic motion when the separation is medium or large ($5.0 < \Delta z < 11.0$ fm), however it is weakly dependent on ϵ when Δz is small ($2.0 < \Delta z < 5.0$ fm).

3 Classically dynamical behaviour of a nucleon in heavy nuclei

For a nucleon, its corresponding classical dynamics is governed by a set of canonical equations of Hamiltonian

$$H = \frac{p^2}{2m_0} + V(\rho, z).$$

As is well known, Poincaré section is a convenient and reliable criterion to judge the regularity (or chaoticity) of a classical system. In the present work the surface of section is chosen to be the plane z-P_z at ρ=0, in which the z-component of angular momentum is a constant (equal to zero). It has been noticed that for other surfaces of section, their qualitative behaviour is similar to that of the section z-P_z. The numerical calculations in the present work are carried out for a nucleon in $^{238}_{92}U$. For other heavy nuclei, the dynamical behaviour is qualitatively the same as that for $^{238}_{92}U$.

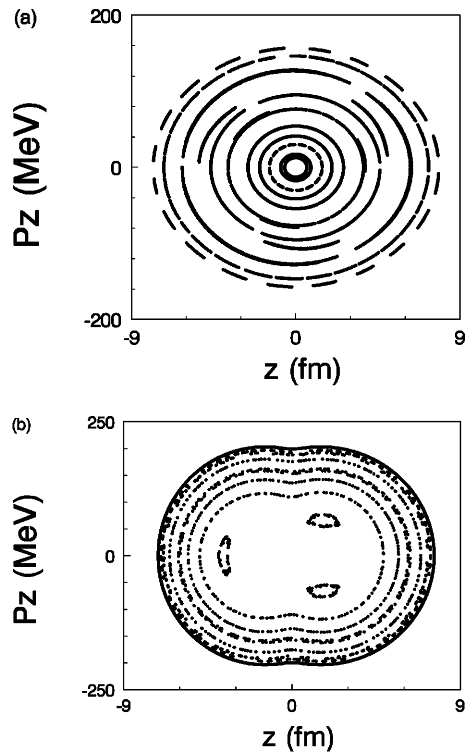

Fig. 1. The Poincare surface of section $z - P_z$ for a nucleon in $^{238}_{92}U$. Δz=3.0 fm, $\beta_1 = 1.0 = \beta_2$, X_i=0, the nucleonic energy e=40.0 MeV. **a** $\epsilon = 0$, **b** ϵ=0.50

The surface of the section are studied systematically with the change of the shape parameters. When the shape deformation is purely quadrupole only ellipsoidally regular invariant tori appear in the surface for any possible values of the separation as shown in Fig. 1a where Δz=3.0 fm, ϵ=0, X_i=0 and $\beta_1 = 1 = \beta_2$. Once a neck deformation is added to a purely quadrupole deformation (the deformation becomes quadrupole-like one), it is found that deformed ellipsoidally invariant tori are dominant and some islands emerge in the surface. We illustrate this situation in Fig. 1b where Δz=3.0 fm, ϵ=0.50, X_i=0 and $\beta_1 = 1 = \beta_2$, from which one can see that there are three moon-shape islands besides the numerous ellipsoidally deformed invariant tori. Thus it can be concluded that the nucleonic classical motion is regular if the deformation is quadrupole or quadrupole-like.

We find that an octupole-like deformatiom strongly influences the classical dynamics. If a small octupole-like is added to a quadrupole or quadrupole-like deformation, a chaotic sea emerges in a small part of the surface, and de-

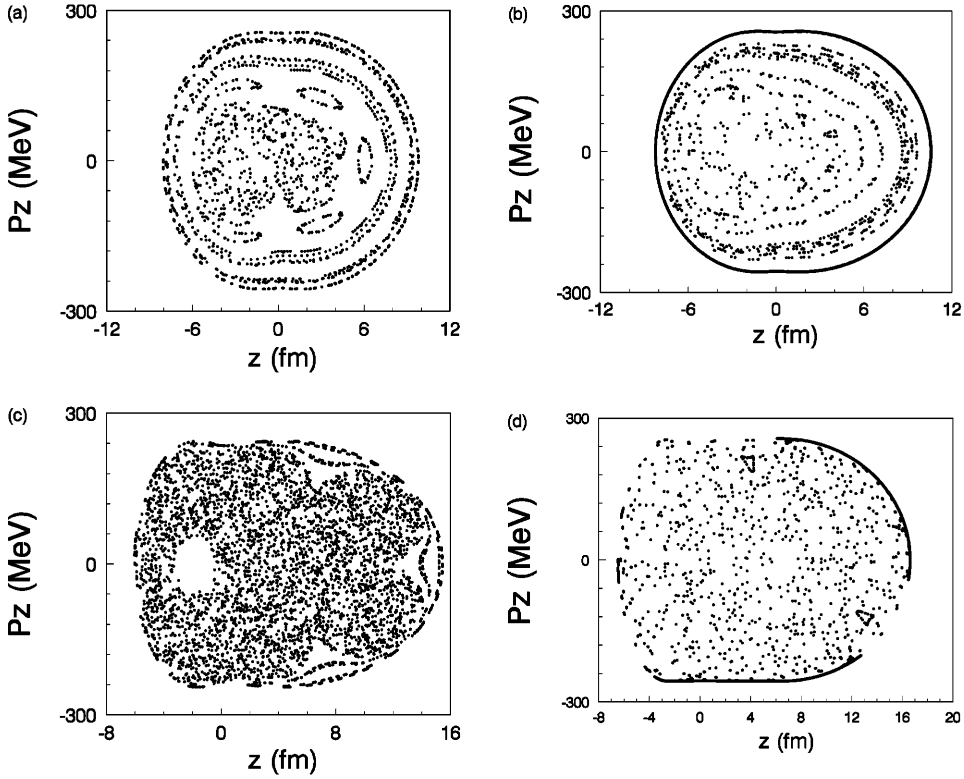

Fig. 2. The Poincare surface of section $z - P_z$ for a nucleon in $^{238}_{92}U$ for different octupole-like deformations. e=40.0 MeV. **a** Δz=3.0 fm, $\beta_1 = 0.80$, $\beta_2 = 1.20$, X_i=0.15, $\epsilon = 0.50$; **b** Δz=3.0 fm, $\beta_1 = 0.70$, $\beta_2 = 1.50$, X_i=0.25, $\epsilon = 0.5$; **c** Δz=3.0 fm, $\beta_1 = 0.40$, $\beta_2 = 2.50$, X_i=0.40, $\epsilon = 0.5$; **d** Δz=9.0 fm, $\beta_1 = 0.4$, $\beta_2 = 2.50$, X_i=0.4, $\epsilon = 0.5$

formed invariant tori and islands appear in the rest of the surface. This situation is illustrated in Fig. 2a in which we take X_i=0.15, $\beta_1 = 0.80$, $\beta_2 = 1.20$ (indicating a small octupole-like deformation) and Δz=3.0 fm, ϵ=0.30. With the increase of octupole-like deformation, more and more invariant tori are damaged and a chaotic sea can be expected to become larger and larger in the surface. In Fig. 2b, X_i=0.25, $\beta_1 = 0.70$, $\beta_2 = 1.50$ (which is obviously a large octupole-like deformation compared to that in Fig. 2a), and the separation and neck deformation are the same as those in Fig. 2a. It can be seen that the chaotic sea in Fig. 2b is apparently larger than that of Fig. 2a. The dynamics in Figs. 2a–b is in between regularity and chaoticity. If an octupole-like deformation becomes so large that the values of the five shape parameters are located in the chaotic region of shape parameter, then most of tori disappear and a chaotic sea takes over the surface. Thus, the classical motion is chaotic. This is exemplified both by Fig. 2c for the part (I) and by Fig. 2d for the part (II) of the chaotic region. Therefore the occurrence of the classically chaotic dynamics requires a pronounced octupole-like deformation.

It is noticed that for small separations the chaotic dynamics can be produced even if the neck deformation is very small. In Fig. 3a the values of shape parameter are taken the same as those in Fig. 2c except for the neck deformation, ϵ=0.02. One can see that in Fig. 3a the dynamics is also chaotic. However, when the separation becomes medium or large the occurrence of chaotic dynamics requires a considerable neck deformation. In Fig. 3b the same values of shape parameter as those in Fig. 2d are taken except for the neck deformation ϵ=0.15. It is shown that the dynamics is in between regularity and chaoticity and obviously different from that in Fig. 2d. We find that a very large neck deformation can suppress the chaotic dynamics especially for medium or large separations. In Fig. 3c we take the same values of shape parameter as those in Fig. 2d except for the neck deformation ϵ=0.95. One can observe that invariant tori are dominant, the motion is approximately regular.

It is shown that for prolate nuclei (corresponding to medium or large separations the chaotic dynamics only occurs if the nucleonic energy is higher than a certain value. In Figs. 4a-b, the values of shape parameter are the same as those in Fig. 2d, but the energy e is taken as 20.0 MeV and 9.0 MeV respectively. From Figs. 4a–b together with Fig. 2d, one can find that with the decrease of energy, the dynamical pattern is changed from chaoticity to regularity via a mixture of regularity and chaoticity. Nevertheless, for oblate nuclei (corresponding to small separations), the situation is quite different. In Figs. 4c–d, the shape parameters are the same values as those in Fig. 2c but the energy e is 3.0 MeV and 2.0 MeV respectively. Fig. 4c shows that when e is as low as 3.0 MeV the dynamics is nearly chaotic. Even if e is reduced to 2.0 MeV the dynamics still shows a mixture of regularity and chaoticity.

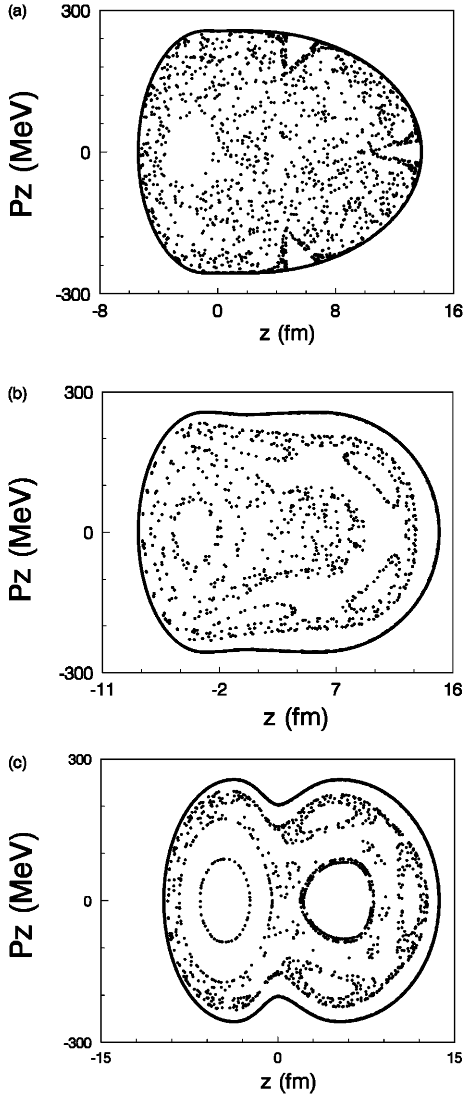

Fig. 3. The Poincare surface of section $z - P_z$ for a nucleon in $^{238}_{92}U$ for different neck deformations. $\beta_1 = 0.40$, $\beta_2 = 2.50$, X_i=0.40, e=40.0 MeV. **a** Δz=3.0 fm, $\epsilon = 0.02$; **b** Δz=9.0 fm, $\epsilon = 0.15$; **c** Δz=9.0 fm, $\epsilon = 0.95$

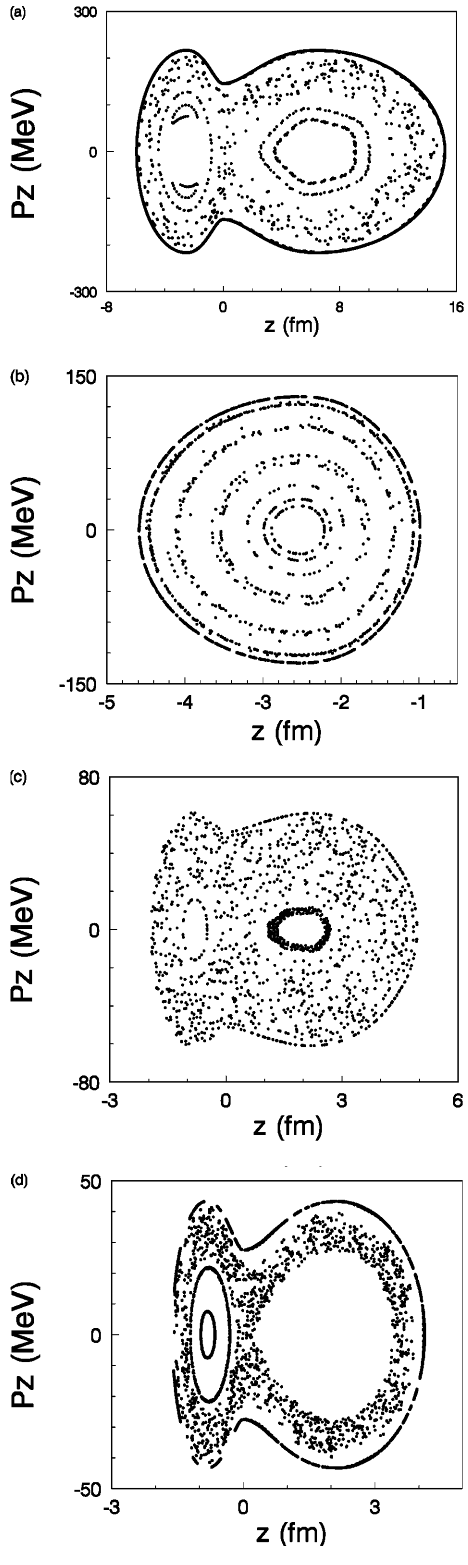

Fig. 4. The Poincare surface of section $z - P_z$ for a nucleon in $^{238}_{92}U$ for different nucleonic energies. $\beta_1 = 0.40, \beta_2 = 2.50$, X_i=0.40, $\epsilon = 0.50$. **a** Δz=9.0 fm, e=20.0 MeV; **b** Δz=9.0 fm, e=9.0 MeV; **c** Δz=3.0 fm, e=3.0 MeV; **d** Δz=3.0 fm, e=2.0 MeV

The above phenomenon can be understood by means of the TCSM potential described by (6). The potential as a function of coordinates is drawn in Fig. 5a for a prolate nucleus and in Fig. 5b for an oblate one. In Fig. 5a and Fig. 5b, the shape parameters have the same values as those in Fig. 2d and in Fig. 2c separately. It is shown that there is a pronounced potential barrier at the position z=0 for the prolate case, however, there is a very low potential barrier at z=0 for the oblate case. Certainly, as displayed in Fig. 5c the height of the barrier is determined by the neck deformation ϵ. The larger the ϵ is, the higher the barrier. In Fig. 5c it can be seen that for the oblate case, the height of barrier is still very low even if ϵ=0.90. On the contrary, for prolate case the height is as high as 25 MeV at ϵ=0.90. As a usual, the potential barrier promotes the nucleonic chaotic motion due to its scattering to the nucleon. This is the reason why the chaotic motion occurs for prolate

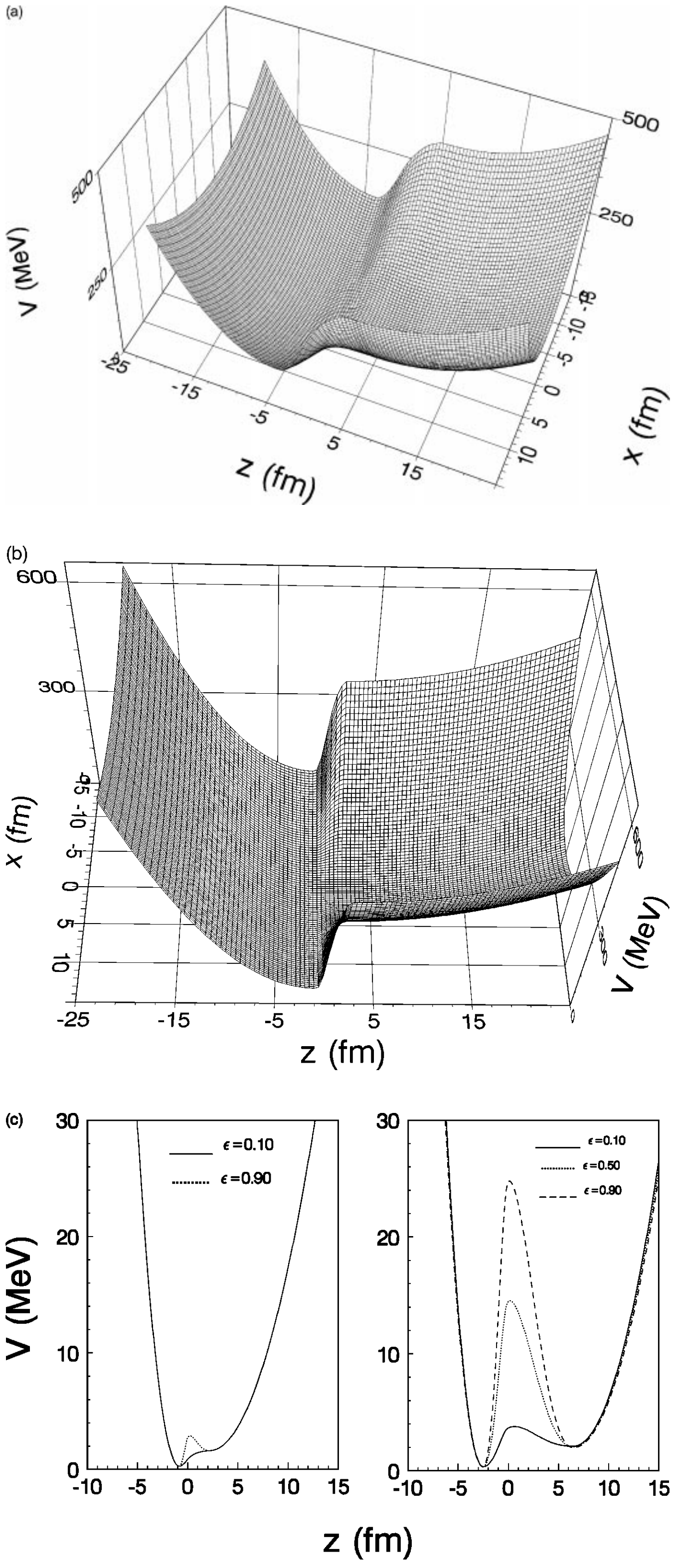

Fig. 5. The potential versus coordinates for prolate and oblate nuclei in the presence of a pronounced octupole-like deformation. $\beta_1 = 0.40, \beta_2 = 2.50$, X_i=0.40. **a** Δz=9.0 fm, $\epsilon = 0.50$; **b** Δz=3.0 fm, $\epsilon = 0.50$; **c** left for Δz=3.0 fm, right for Δz=9.0 fm

nuclei with a considerable neck deformation, and why no chaotic motion is expected for prolate nuclei in the case of Heiss et al. (neck deformation was not included). However, if the nucleonic energy approaches to (but above) the height of the barrier, the nucleonic kinetic energy becomes quite small once the nucleon is nearby the top of the barrier, which makes the nucleonic trajectories stable. Then, the chaotic motion is suppressed. If the energy is below the height, the nucleon moves within a parabolic-like potential, a regular or mixed dynamics can be expected. Therefore, due to a pronounced barrier in the prolate case, the chaotic dynamics is only produced if the nucleonic energy has a considerable value (apparently higher than the height of the potential barrier). Nevertheless, due to a low barrier in the oblate case, the chaotic dynamics appears even if the energy is very low.

Comparing the nucleonic quantum motion with its corresponding classical dynamics, one can find that for a nucleon in heavy nuclei, there is a good quantum-classical correspondence of regular (chaotic) motion as long as the energy of a nucleon is obviously higher than the height of the barrier, so that BGS conjecture [4] is corroborated once again here.

Recently, Wilkinson [18] pointed out that the adiabatic time-dependent chaotic Hamiltonians can be a source of dissipation, and the dissipation is drastically reduced when the Hamiltonian is integrable. Similar conclusion was drawn by Carvalho et al and by Blocki et al. [19]. In a nucleus, there exist collective and intrinsic (nucleonic) motions. The former which leads to the change of nuclear shape can be regarded as slow variables, the latter rapid variables. This is so-called adiabatic approximation. In the past several decades the induced heavy nuclei fission have been widely investigated [20–22]. During the fission process, strong nuclear dissipation has been found for a long time. The present study has shown that the chaotic motion of a nucleon in heavy nuclei is strongly dependent on the nuclear shape deformations. From the pointview of Wilkinson, the nuclear disspation of heavy nuclei should be the shape dependent, and strong nuclear dissipation is expected for the nuclei with the shape parameter values described by the part (II) of the chaotic region (An oblate nuclear shape is impossible during the fission process). The present work provides us a dynamical understanding of the shape dependence of nuclear dissipation.

4 Conclusions

In the present paper we have studied the classical dynamics of a nucleon in heavy nuclei within the framework of the TCSM when shape parameters are varied systematically. It has been found that a pronounced octupole-like deformation is necessary to produce the classically chaotic dynamics. The effect of neck deformation on the classical dynamics for a small separation is quite different from that for a medium or large separation. In the case of medium or large separations, a considerable neck deformation and a higher energy value over the height of the barrier are required for the occurrence of the classical chaos. It has been found that when the energy is apparently higher than the height of the barrier there is a good quantum-classical correspondence of nucleonic regular (chaotic) motion by comparing the classically dynamical behaviour with the statistical properties of the quantum spectrum. So that the BGS conjecture has been confirmed once again for such a complex system. It has been shown that even if a nucleus appears to be prolate one, the classical chaos can also occur provided that there exist a considerable neck and a pronounced octupole-like deformation. This is quite different from the conclusion of Heiss [11] et al. It is suggested that the difference between the potential barrier for prolate nuclei and that for oblate ones is reponsible for the energy-dependence difference between the nucleonic chaotic dynamics for prolate nuclei and that for oblate ones. In addition, it is implied that the nuclear dissipation should be shape-dependent, and strong nuclear dissipation is expected for the nuclei with the shape parameter values described by the part (II) of the chaotic region. It provides us a dynamical understanding of shape dependence of nuclear dissipation. The present work is heuristic to the study of a finite system where pure geometry dominates the dynamical behavior of the system.

We would like to express our gratitute to Prof. Jun-qing Li for useful discussions and suggestions. This work is supported by the National Natural Science Foundation of China and China Postdoctoral Science Foundation.

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