Classical Dynamics in Deformed Spaces

A. Leznov and J. Mostovoy

Abstract

We consider a 3-parametric linear deformation of the Poisson brackets in classical mechanics. This deformation can be thought of as the classical limit of dynamics in so-called “quantized spaces”. Our main result is a description of the motion of a particle in the corresponding Kepler-Coulomb problem.

1 Introduction

The set-up of a problem in quantum mechanics usually consists of specifying (a) the structure of the algebra of the observables and (b) the Hamiltonian of the problem. If the commutation relations in the algebra of the observables are deformations of the usual commutation relations between the coordinates and the momenta, the situation is often referred to as ”dynamics in a quantized space”. Schrödinger was apparently the first to consider such a situation: in [1] he calculated the energy levels for the Coulomb potential on a 3-sphere. For this problem the term ”quantized space” is rather appropriate as the spectrum of the distance operator on a 3-sphere is discrete even in the absence of any field.

One can obtain ”classical dynamics in a quantized space” by setting

$$\lim_{\hbar \to 0} \frac{i}{\hbar}[A, B] = \{A, B\}$$

where \{·,·\} are the Poisson brackets of classical dynamics. We prefer to speak about ”classical dynamics in deformed spaces” (keeping in mind that it is not just the space but the whole dynamics being deformed).

In fact, one does not need a reference to the quantum theory in order to construct a deformation of the functional algebra of classical observables. However, we are specifically interested in the classical version of the 3-parametric family of deformations considered in [2, 3, 4] in the context of quantum mechanics.

This family arises naturally if one considers deformations of the commutation relations involving the four space-time coordinates, ten generators of the Poincare algebra (that is, infinitesimal translations and proper Lorentz transformations) and the unity element, with the property that the resulting algebra contains the Lorentz algebra as a subalgebra.

In what follows we will consider only three-dimensional non-relativistic problems. The usual commutation relations between the dynamic variables are replaced by the following
set of brackets satisfying the Jacobi identity:

\[
\begin{align*}
\{p_i, x_j\} &= \delta_{ij} I - \frac{\epsilon_{ijk} l_k}{S}, \\
\{x_i, x_j\} &= -\frac{\epsilon_{ijk} l_k}{M^2}, \\
\{l_i, x_j\} &= -\frac{\epsilon_{ijk} l_k}{L^2}, \\
\{x_i, p_j\} &= -\frac{\epsilon_{ijk} l_k}{M^2}, \\
\{l_i, p_j\} &= -\frac{\epsilon_{ijk} l_k}{L^2}, \\
\{I, x_i\} &= \frac{x_i}{S} - \frac{p_i}{M^2}, \\
\{I, p_i\} &= \frac{x_i}{L^2} - \frac{p_i}{S}, \\
\{I, l_i\} &= 0.
\end{align*}
\] (1)

Here \(x_i\) are the coordinates and \(p_i\) and \(l_i\) are the components of the momentum and angular momentum respectively; the function \(I\) is a deformation of the unity function. The three parameters \(L^2, M^2\) and \(S\) have the dimensions of area, square of the momentum and action respectively.

In general, these parameters may depend on rotation-invariant functions of the dynamic variables such as \(x^2, l^2, (px)\) etc. In this case in the resulting non-linear functional algebra the Jacobi identities are still satisfied. Here, however, we only consider the case of \(L^2, M^2\) and \(S\) being constants. It should be stressed that we do not require the signs of \(L^2\) and \(M^2\) to be positive.

In order to obtain the flat space with the usual dynamics on the scale of the Solar system one has to assume that the values of \(L^2, M^2\) and \(S\) are huge. In fact, taking \(M^2\) and \(S\) to be infinite, we get the dynamics in the space of constant curvature equal to \(1/L^2\), so \(L\) is on the scale of the size of the Universe.

An important feature of the relations (1) is their non-invariance with respect to the time reversal. Namely, changing the signs of \(p, l\) and of the brackets simultaneously one obtains relations of the same form but with the sign of the parameter \(S\) reversed.

Various particular cases of the above commutation relations are studied in the literature. The case of the space of constant curvature (that is, \(M^2, S = \infty\)) has received most attention. The first work related to this case is due to Schrödinger [1] (who assumed the curvature to be positive, for the case of negative curvature see [5]). Many other works on the subject have appeared since; we only mention [6] which contains a discussion of the classical mechanics (and, in particular, the Kepler problem) in the space of constant curvature.

The case \(S = \infty\) is due to Yang [8]. The general case of the above commutation relations was introduced in [2], see also [4]. The quantum energy levels for the Kepler-Coulomb potential and the harmonic oscillator for this general case are calculated in [3].

Other approaches to deformed spaces exist: see, for example, [9].

About the notation: we use bold italic letters for vectors. The scalar and vector products of vectors \(a\) and \(b\) will be denoted by \((ab)\) and \([ab]\) respectively, the absolute value of \(a\) is denoted by \(|a|\) and its square - by \(a^2\).

For simplicity, we always assume the particles to be of unit mass.

## 2 The canonical coordinates

### 2.1 The canonical coordinates

Depending on the signs of the constants \(L^2, M^2\) and \(L^2M^2 - S^2\), the functional algebra (1) is equivalent as a Lie algebra to one of the algebras \(o(5), o(1, 4)\) or \(o(2, 3)\) ([3]). According to the Darboux theorem these algebras can be resolved in terms of four pairs of independent canonically conjugate coordinates and momenta and two cyclic variables constructed from
the elements of (1). In the algebraic language, the cyclic variables are the Casimir operators of the corresponding algebra.

The Casimir operators of the algebra under consideration are as follows:

\[ K_2 = I^2 + \frac{x^2}{L^2} + \frac{p^2}{M^2} - \frac{2(px)}{S} - l^2 \left( \frac{1}{S^2} - \frac{1}{L^2M^2} \right) \]

and

\[ K_4 = (I - [x, p])^2 - \frac{(pl)^2}{M^2} + \frac{2(pl)(xl)}{S} - \frac{(xl)^2}{L^2}. \]

The Poisson brackets of \( K_2 \) and \( K_4 \) with any function vanish and their values can be fixed. In order to get a correct limit in the usual (that is, nondeformed) case we shall set \( K_2 = 1 \).

As for the value of \( K_4 \), we will only consider the scalar representation of the algebra of the observables, that is, set \( I_l = [xp] \).

This implies that \( (pl) = (xl) = 0 \) and, hence, that \( K_4 = 0 \). The relation (2) generalises the definition of the angular momentum as a function of other variables to the case of the deformed space.

Imposing the relation (2) we reduce the number of Darboux pairs by one. An explicit resolution of the algebra (1) in terms of three angular variables can be found in the Appendix.

2.2 Problems with spherical symmetry

For problems with spherical symmetry (such as the Kepler problem considered below) it may be of use considering the algebra of observables invariant under rotations. Namely, consider the algebra generated by the 5 spherically invariant variables: \( r^2 \equiv x^2 + \frac{l^2}{M^2}, \ p'^2 \equiv p^2 + \frac{l^2}{L^2}, \ D \equiv (px) + \frac{l^2}{S}, \ I, \ l^2 \). The commutation relations of the nonlinear algebra generated by these 5 variables follow directly from (1):

\[ \{I, p'^2\} = 2 \left( \frac{p}{M} - \frac{l^2}{S} \right), \quad \{I, r^2\} = - \left( \frac{p}{M} - \frac{l^2}{S} \right), \quad \{I, D\} = \frac{r^2}{L^2} - \frac{p'^2}{L^2}, \]

\[ \{p'^2, r^2\} = 4DI, \quad \{D, p'^2\} = -2p'^2I, \quad \{D, r^2\} = 2r^2I \] (3)

The nonlinear functional algebra (3) is generated by 5 elements and has three cyclic variables: two of them, namely \( K_2 \) and \( K_4 \), come from the initial algebra. The third cyclic variable is \( l^2 \) which commutes with all the generators of (3). Thus, by the Darboux theorem, this algebra can be resolved in terms of one coordinate and one momentum. We choose \( r \) as the coordinate. (In the non-deformed case this is just the radial variable.)
The values of $K_2$ and $K_4$ being fixed, the quantities $D$, $p'$ and $I$ depend only on $r$, the corresponding canonically conjugate momentum which we denote by $\rho$, and $l^2$. Taking the squares of the vectors on both sides of (2) and using the fact that $K_2 = 1$ after some manipulations we obtain

$$p'^2 = \frac{1}{r^2} \left(D^2 + l^2\right). \quad (4)$$

From the last equality of (3) we get that

$$I = \frac{1}{r} D_\rho$$

Substituting this expression together with (4) into the equality $K_2 = 1$ we obtain an equation on $D$

$$1 = \frac{D^2}{r^2} + \frac{D^2 + l^2}{M^2 r^2} + \frac{r^2}{L^2} - \frac{2D}{S} + l^2 \left(\frac{1}{S^2} - \frac{1}{L^2 M^2}\right)$$

with the general solution

$$D = \frac{M^2 r^2}{S^2} - \sqrt{\left(1 + r^2 M^2 \left(\frac{1}{S^2} - \frac{1}{L^2 M^2}\right)\right) \left(M^2 r^2 - l^2\right) \cos \frac{\rho}{M}}. \quad (5)$$

As a corollary we have

$$I = \sqrt{\left(1 + r^2 M^2 \left(\frac{1}{S^2} - \frac{1}{L^2 M^2}\right)\right) \left(1 - \frac{l^2}{M^2 r^2}\right) \sin \frac{\rho}{M}} \quad (6)$$

Thus all elements of the nonlinear functional algebra given by (3) are expressed in terms of the canonically conjugate pair $(r, \rho)$ and the square of the angular momentum $l^2$.

## 3 Free motion

The Hamiltonian of the free motion is chosen so as to commute with infinitesimal translations and rotations, that is, with $p$ and $l$. The simplest such expression which gives the correct limit as the parameters tend to infinity is

$$E = \frac{1}{2} \left(p^2 + \frac{l^2}{L^2}\right).$$

It follows immediately from the commutation relations that the vectors $p$ and $l$ are constant. Also,

$$\dot{x} = l p - \frac{[xl]}{L^2} + \frac{[pl]}{S}$$

and

$$\dot{I} = \frac{\rho^2}{S} - \frac{(px)}{L^2}. \quad (4)$$
From this

\[ \dot{I} + \frac{2E}{L^2} I = 0 \]

and, hence,

\[ I = I_0 \cos \left( \frac{\sqrt{2E}}{L} (t - t_0) \right). \]

Now, taking the vector product of (2) with \( p \) one has

\[ (px)p - p^2 x = I[p] \]

from where

\[ x = \left( \frac{L^2}{S} + \frac{I_0}{p^2} L \sqrt{2E} \sin \left( \frac{\sqrt{2E}}{L} (t - t_0) \right) \right) p - \frac{I_0}{p^2} \cos \left( \frac{\sqrt{2E}}{L} (t - t_0) \right) [lp]. \] (7)

The value of \( I_0 \) can be found from the condition \( K_2 = 1 \):

\[ I_0 = \frac{p}{\sqrt{2E}} \sqrt{1 + 2E \left( \frac{L^2}{S^2} - \frac{1}{M^2} \right)}. \]

In the case \( L^2 = \infty \) the expression for \( x \) should be modified to read

\[ x = x_0 + \left( I_0 p + \frac{[pl]}{S} \right) t + \frac{p^2}{2S} pt^2 \]

with

\[ I_0 = \sqrt{1 - \frac{p^2}{M^2} + \frac{l^2}{S^2} + \frac{2(px_0)}{S}}. \]

The orbits of the free motion are elliptic, parabolic or hyperbolic according to whether \( L^2 \) is positive, zero, or negative. This should not be surprising given that \( L^{-2} \) can be interpreted as the curvature of the space. For example, set \( M^2, S = \infty \) and take \( L^2 > 0 \). Then the elliptic orbits (7) can be obtained from the geodesics on the 3-sphere

\[ 1 = I^2 + \frac{x^2}{L^2} \]

by projecting them orthogonally on a Euclidean 3-space. In particular, the farthest distance from the origin reached by a freely moving particle is the same for all orbits and equals to \( L \).
4 The Kepler problem

4.1 The equation of the orbit

The Hamiltonian for the Kepler problem is chosen (as in [3]) so as to commute with the R"unge-Lenz vector $A$ defined as

$$A = [pl] + \frac{\alpha x}{r} = \frac{1}{I} \left( p^2 x - (px)p \right) + \frac{\alpha x}{r}$$

(8)

where $r$ is given by $r^2 = x^2 + \frac{l^2}{M^2}$.

The commutation relations for the components of $A$ are as follows:

$$\{A_i, A_j\} = 2 \left( E - \frac{l^2}{L^2} \right) \epsilon_{ijk} l_k,$$

where $E$ stands for the expression

$$E = \frac{1}{2} \left( p^2 + \frac{l^2}{L^2} + \frac{2\alpha I}{r} - \frac{\alpha^2}{M^2 r^2} \right).$$

(9)

Also,

$$\{l_i, A_j\} = -\epsilon_{ijk} A_k$$

and, hence,

$$\{l^2, A_j\} = 2\epsilon_{ijk} l_j A_k.$$

To see that $E$ commutes with all of the components of $A$ first notice that

$$A^2 = \alpha^2 + l^2 \left( 2E - \frac{l^2}{L^2} \right).$$

(10)

Now,

$$2\{E, A_i\} = \left\{ \frac{A^2}{l^2}, A_i \right\} - \alpha^2 \left\{ \frac{1}{l^2}, A_i \right\} + \frac{1}{L^2} \left\{ l^2, A_i \right\} = 0.$$

Thus $E$ can be chosen to be the Hamiltonian for the Kepler problem.

The three vectors $x$, $p$ and $A$ lie in the plane orthogonal to the vector of angular momentum. Let us choose the polar coordinates $(|x|, \phi)$ in this plane with the direction $\phi = 0$ being given by the R"unge-Lenz vector. (If $A = 0$ we take $\phi = 0$ to be an arbitrary ray passing through the origin.)

Taking the scalar product of (8) with $x$ one has

$$\frac{A|x| \cos \phi - \alpha r}{l^2} = I - \frac{\alpha}{M^2 r}.$$

(11)

Taking the vector product of $A$ and $l$ we have

$$[Al] = -l^2 p + \frac{\alpha}{r} [xl].$$

(12)
The scalar product of (12) with $x$ gives

$$(px) = -\frac{1}{l^3}([Al], x) = \frac{A|x| \sin \phi}{l}.$$  

(13)

Now the orbits can be found as follows. In the expression $K_2 = 1$ one can substitute $(px)$ from (13), $I$ — from (11) and $p^2$ — from (9) and (10). The resulting equation of the orbit is as follows:

$$1 + 2\frac{A|x| \sin \phi}{lS} = \frac{(A|x| \cos \phi - \alpha r)^2}{l^4} + \frac{r^2}{L^2} + \frac{A^2 - \alpha^2}{l^2 M^2} - \frac{l^2}{S^2}. \tag{14}$$

In the Cartesian coordinates $(x, y) = (|x| \cos \phi, |x| \sin \phi)$ the equation (14) takes the following form:

$$x^2 \left( \frac{A^2 + \alpha^2}{l^4} + \frac{1}{L^2} \right) + y^2 \left( \frac{\alpha^2}{l^4} + \frac{1}{L^2} \right) - \frac{2A \alpha}{l^4} \sqrt{\frac{x^2}{l^2} + \frac{y^2}{M^2}} - l^2 \left( \frac{1}{S^2} - \frac{1}{L^2 M^2} \right) - 1 = 2\alpha A \frac{l^4}{l^4} \sqrt{x^2 + y^2 + \frac{l^2}{M^2}}. \tag{15}$$

This equation defines a branch of a plane quartic. The other branch of this quartic describes the orbit with the same values of $A$ and $l$ but with the opposite value of $\alpha$. Notice that changing $l$ for $-l$ we obtain, generally, a different orbit. (Recall that the commutation relations (1) are not invariant under time reversal.)

Interestingly, it was Giovanni Domenico Cassini (1625–1712) who first suggested that the trajectories of planets are quartic curves. Now these curves are known as Cassinian ovals. A Cassinian oval is defined as the locus of points with the product of their distances from two fixed points equal to a fixed constant.

4.2 The parametric form for the orbit

From now on we will assume for simplicity that $L^2$ is positive. Taking into the account the fact that $K_2 = 1$, this implies, in particular, that all orbits are closed.

The curve (15) can be parametrised as follows. The condition $L^2 > 0$ means that we can choose $a, b$ so that

$$a^2 + b^2 = \frac{\alpha^2 + A^2}{l^4} + \frac{1}{L^2}$$

and

$$ab = \frac{\alpha A}{l^4}.$$

Explicitly,

$$a, b = \frac{1}{2} \left( \frac{\alpha^2 + A^2}{l^4} + \frac{1}{L^2} \right) \pm \sqrt{\left( \frac{\alpha^2 + A^2}{l^4} + \frac{1}{L^2} \right)^2 - 4\frac{A^2 \alpha^2}{l^8}}.$$
and we take \( a^2 > b^2 \).

The equation (15) can be written as

\[
(ax - br)^2 + \left( \nu y - \frac{A}{lS\nu} \right)^2 = \gamma^2.
\]

Here

\[
\nu^2 = \frac{1}{2} \left( \frac{a^2 - A^2}{l^4} + \frac{1}{L^2} + \sqrt{\left( \frac{a^2 + A^2}{l^4} + \frac{1}{L^2} \right)^2 - \frac{4A^2\alpha^2}{l^4}} \right)
\]

and

\[
\gamma^2 = 1 + \nu^2 \left( \frac{1}{S^2} - \frac{1}{L^2M^2} \right) \frac{a^2}{\nu^2}.
\]

In the non-deformed limit \( \gamma = 1 \) and

\[
\nu^2 = \frac{\alpha^2 - A^2}{l^4} = -\frac{2E}{l^2}
\]

for \( E < 0 \) and \( \nu^2 = 0 \) otherwise.

We can set

\[
ax - br = \gamma \cos \theta,
\]

\[
\nu y - \frac{A}{lS\nu} = \gamma \sin \theta.
\]

so that

\[
x = \frac{1}{a^2 - b^2} \left( a \gamma \cos \theta + b \frac{|a|}{a} \sqrt{\nu^2 \cos^2 \theta + (a^2 - b^2) \left( \frac{\gamma \sin \theta + \frac{A}{lS\nu^2}}{\nu^2 + \frac{1}{M^2}} \right)^2} \right),
\]

\[
y = \gamma \nu \sin \theta + \frac{A}{lS\nu^2}.
\]

In the classical limit the above parametrisation gives the usual trigonometric parametrisation of the elliptic trajectory. There exists also a generalisation of the hyperbolic parametrisation.

Write (15) as

\[
(bx - ar)^2 - \left( \mu y + \frac{A}{lS\mu} \right)^2 = \delta^2
\]

with

\[
\mu^2 = -\frac{1}{2} \left( \frac{a^2 - A^2}{l^4} + \frac{1}{L^2} \right) \left( \frac{a^2 + A^2}{l^4} + \frac{1}{L^2} \right)^2 - \frac{4A^2\alpha^2}{l^4}
\]

(17)

8
and
\[ \delta^2 = 1 - l^2 \left( \frac{1}{S^2} - \frac{1}{L^2 M^2} \right) \frac{b^2}{\mu^2}. \]

Assume that \( \delta^2 > 0 \). Then setting
\[ bx - ar = \delta \cosh \theta, \]
\[ \mu y + \frac{A}{l S \mu} = \delta \sinh \theta. \]

we obtain the hyperbolic parametrisation of the orbit. (The case \( \delta^2 < 0 \) is entirely similar.)

Notice that \( \mu^2 \nu^2 = A^2 l^4 \), so in the non-deformed limit only one of the parametrisations is meaningful for the given values of \( \alpha \) and \( A \), the other just giving the equation of the orbit.

In the deformed case both parametrisations can be meaningful for the same closed orbit.

In this case the hyperbolic parametrisation is two-valued and the parameter \( \theta \) belongs to a closed interval.

In the non-deformed limit the hyperbolic parametrisation remains two-valued with the parameter \( \theta \) varying over all real numbers. The two branches correspond to the two possible values of \( I = \pm 1 \). One branch gives the usual parametrisation of a hyperbolic trajectory. The second branch parametrises a trajectory with the same value of \( A \) and \( l \) but with the opposite sign of \( \alpha \).

### 4.3 Dependence on time

Taking the Poisson bracket of \( x \) with the Hamiltonian we obtain the equations of motion:
\[ \dot{x} = lp - \frac{[x]}{L^2} + \frac{[pl]}{S} + \frac{\alpha x}{S r} - \frac{\alpha p}{M^2 r}. \]

Substituting \( I \) and \( p \) from (11) and (12) respectively, we get
\[ \dot{x} = \frac{1}{l^4} (Ax - \alpha r) \left( \frac{\alpha}{r} [x] - [Al] \right) - \frac{[x]}{L^2} + \frac{A}{S}. \]

In particular, for the coordinate \( y \) we have:
\[ \dot{y} = \frac{1}{l^4} \left( \left( A^2 + \alpha^2 + \frac{l^4}{L^2} \right) x - \alpha A \left( \frac{x^2}{r} + r \right) \right), \]
or, equivalently,
\[ \dot{y} = -\frac{l}{r} (ax - br)(bx - ar). \]

It follows from (16) that
\[ \frac{r}{bx - ar} = -\frac{|a|}{a} \cdot \left( \frac{br \cos \theta}{\sqrt{\Delta}} + |a| \right) \cdot \frac{1}{a^2 - b^2}. \]
where $\Delta$ is the expression under the square root in (16). Thus
\[
\frac{|a|}{a} \cdot b \nu(a^2 - b^2) \cdot dt = d\theta \left( |a| + \frac{b \gamma \cos \theta}{\sqrt{\Delta}} \right).
\] (18)

The function $\Delta$ only depends on $\sin \theta$ so the period of the motion can be found without integrating (18):

\[
T = \frac{2\pi|a|}{b \nu(a^2 - b^2)}.
\]

Notice that for $A, l$ and $\alpha$ fixed, the period only depends on $L^2$ and not on $M^2$ and $S$.

One can re-write (18) as
\[
\frac{|a|}{a} \cdot (a^2 - b^2) dt = \frac{|a|}{\nu} d\theta + \frac{b}{\mu} \sqrt{q^2 + \left(1 - \frac{1}{S^2} \frac{1}{M^2} \frac{1}{L^2} \right) \left(\nu^2 L^2 - 1\right)}
\]
where
\[
q = \frac{\mu}{\nu} \gamma \sin \theta + \frac{(a^2 - b^2) l L}{\nu S}
\]
and $\mu$ is defined as in (17) so that $\mu^2 + \nu^2 = a^2 - b^2$. Integrating, one gets the following parametrisation of the time $t$ by $\theta$:

\[
\frac{|a|}{a} \cdot (t - t_0)(a^2 - b^2) l = \frac{|a|}{\nu} \theta + \frac{b}{\mu} \sinh^{-1} \left( \frac{\frac{\mu}{\nu} \gamma \sin \theta + \frac{(a^2 - b^2) l L}{\nu S}}{\sqrt{1 + \frac{1}{L^2} \left(\frac{1}{S^2} - \frac{1}{M^2} \frac{1}{L^2} \right) \left(1 - \nu^2 L^2\right)}} \right).
\] (19)

(The above answer is valid if
\[
1 + \frac{l^2}{L^2} \left(\frac{1}{S^2} - \frac{1}{M^2} \frac{1}{L^2} \right) \left(1 - \nu^2 L^2\right) > 0;
\]
obvious modifications are needed otherwise in (19).)

The standard trigonometric parametrisation of $t$ by $\theta$ in the non-deformed Kepler problem can be obtained from (19) by setting first $S$ and $M^2$ to be infinite and then taking the limit as $L^2$ tends to infinity. (Note that as $L^2$ tends to $\infty$, the parameter $\mu$ tends to 0, as $\mu = A \nu l L$.)

5 The Hamilton-Jacobi equation

The spherically symmetric part of the Hamilton-Jacobi equation is obtained by substituting (4), (5) and (6) into the explicit expression for the Hamiltonian function
\[
2E = \frac{1}{\nu^2} (D^2 + l^2) + \frac{2\alpha}{\nu} I - \frac{\alpha^2}{M^2 l^2}.
\]

10
Namely, we obtain
\[ 2E r^2 - l^2 + \frac{\alpha^2}{M^2} = \left( \frac{M^2 r^2}{S} - \sqrt{\left( 1 + r^2 M^2 \left( \frac{1}{S^2} - \frac{1}{L^2 M^2} \right) \right)} \right) (M^2 r^2 - l^2) \cos \frac{W_r}{M}^2 \]
\[ + \frac{2\alpha}{M} \sqrt{\left( 1 + r^2 M^2 \left( \frac{1}{S^2} - \frac{1}{L^2 M^2} \right) \right)} (M^2 r^2 - l^2) \sin \frac{W_r}{M} \]

where \( W \) is the radial part of the action function. One arrives to the same expression considering the quasi-classical limit for the Schrödinger equation in a deformed (quantized) space.

We do not know a direct method of solving this (rather unusual at the first sight) equation. In fact, one only needs to know the the second derivative \( W_{rE} \) in order to describe \( r \) as a function of time: \( r(t) \) can be found from the equality
\[ t - t_0 = \int dr W_{rE}. \]

Below we obtain an algebraic equation for the second derivative \( W_{rE} \) of the radial part of the action.

Let us introduce the following notation:
\[ u = \frac{M^2 r^2}{S} - \sqrt{\left( 1 + r^2 M^2 \left( \frac{1}{S^2} - \frac{1}{L^2 M^2} \right) \right)} \left( M^2 r^2 - l^2 \right) \cos \frac{W_r}{M}, \]
\[ v = \sqrt{\left( 1 + r^2 M^2 \left( \frac{1}{S^2} - \frac{1}{L^2 M^2} \right) \right)} \left( M^2 r^2 - l^2 \right) \sin \frac{W_r}{M} - \frac{\alpha}{M}. \]

In this notation the Hamilton-Jacobi equation becomes
\[ u^2 + \frac{2\alpha}{M} v = 2E r^2 - l^2 - \frac{\alpha^2}{M^2} \equiv \theta. \]

Taking the derivative of both sides with respect to \( E \) we have
\[ uv = Mr^2 \left( \frac{1}{W_{rE}} - \frac{\alpha}{S} \right) \equiv F. \]

Finally, taking into the account that \( \sin^2 x + \cos^2 x = 1 \) we obtain
\[ v^2 - \frac{2M^2 r^2}{S} u = M^2 r^2 \left( 1 - \frac{2E}{M^2} - \frac{r^2}{L^2} - l^2 \left( \frac{1}{S^2} - \frac{1}{M^2 L^2} \right) \right) \equiv \bar{\theta} \]

The condition of self-consistency of the three equations above with respect to only two functions \( u, v \) leads to an equation for the function \( F \) in the form
\[ (\overline{F}^2 - \theta \bar{\theta})^2 = \overline{F} \left( \overline{F}^2 - 3\theta \bar{\theta} \right) + (s^2 \bar{\theta}^3 + w^2 \theta^3) \overline{F} \]
where \( s = \frac{M^2}{s}, w = \frac{M^2 r^2}{S} \) and \( \overline{F} \equiv F + 4sw. \)
6 Conclusion

The main result of the present paper is a parametric solution of the Kepler problem in a deformed space whose construction involves three parameters with dimensions of length, momentum and action respectively.

The orbits of particles in this version of the Kepler problem turn out to be (branches of) quartic curves. This should be compared with the analysis of Higgs [6] who studies the orbits in the case where there is only one non-trivial deformation parameter, namely the length \( L \). In Higgs’ picture the orbits are elliptic, parabolic or hyperbolic just as in the undeformed case. This, however, is due to a special choice of coordinates used in [6]. These coordinates, though well-suited to the geometry of the problem, have no physical significance. The coordinates we use are distinguished among all possible coordinate systems as ”the physical coordinates”. The transformation between Higgs’ coordinates and the physical coordinates is very simple: in our notations, Higgs’ coordinates can be written as \( x_i/I \).

The fact that the Kepler problem can be solved in the deformed space gives some hope that other problems (for example, construction of field theories) could also have reasonable solutions in deformed spaces. The discussion of this subject is, however, beyond the scope of the present article.

Appendix: Structure of the algebra (1).

Throughout the section we assume that \( M^2 > 0 \) and

\[
\frac{1}{S^2} - \frac{1}{L^2 M^2} > 0.
\]

Take the spherical coordinates with the angles \((\phi, \theta, \chi)\) in \( \mathbb{R}^4 \). Let \((u, u_4)\) be the unit vector

\[
(u, u_4) = (\sin \chi \sin \theta \sin \phi, \sin \chi \sin \theta \cos \phi, \sin \chi \cos \theta, \cos \chi)
\]

and set

\[
 p_\chi = \frac{\partial}{\partial \chi},
 p_\theta = \frac{\partial}{\partial \theta},
 p_\phi = \frac{\partial}{\partial \phi}.
\]

Define the vectors \( F \) and \( L \) by

\[
 F_1 = \sin \theta \sin \phi \frac{\partial}{\partial \chi} + \cot \chi \cos \theta \sin \phi \frac{\partial}{\partial \theta} + \cot \chi \frac{\cos \phi}{\sin \phi} \frac{\partial}{\partial \phi},
 F_2 = \sin \theta \cos \phi \frac{\partial}{\partial \chi} + \cot \chi \cos \theta \cos \phi \frac{\partial}{\partial \theta} - \cot \chi \frac{\sin \phi}{\sin \theta} \frac{\partial}{\partial \phi},
 F_3 = \cos \theta \frac{\partial}{\partial \chi} + \cot \chi \sin \theta \frac{\partial}{\partial \phi}
\]

and

\[
 L_1 = \sin \phi \frac{\partial}{\partial \theta} + \cos \phi \cot \theta \frac{\partial}{\partial \phi},
 L_2 = \cos \phi \frac{\partial}{\partial \theta} - \sin \phi \cot \theta \frac{\partial}{\partial \phi},
 L_3 = -\frac{\partial}{\partial \phi}.
\]

respectively.
Now,
\[ \begin{align*}
I &= (uF) + u_4 \rho, \\
\frac{1}{M}(p - \frac{M^2}{S} x) &= \rho u + [Lu] - u_4 F, \\
x \cdot M \sqrt{\frac{1}{S^2} - \frac{1}{L^2 M^2}} &= F, \\
l \cdot \sqrt{\frac{1}{S^2} - \frac{1}{L^2 M^2}} &= L.
\end{align*} \]
\[ (23) \]

Acknowledgments

We would like to thank Carlos Villegas-Blas for useful advice. The second author was supported by the grant J-33616E awarded by CONACyT.

References


A. Leznov: andrey@buzon.uaem.mx
J. Mostovoy: jacob@matcuer.unam.mx