Geometrical origin of the ∗-product in the Fedosov formalism

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Abstract

The construction of the ∗-product proposed by Fedosov is implemented in terms of the theory of fibre bundles. The geometrical origin of the Weyl algebra and the Weyl bundle is shown. Several properties of the product in the Weyl algebra are proved. Symplectic and abelian connections in the Weyl algebra bundle are introduced. Relations between them and the symplectic connection on a phase space $M$ are established. Elements of differential symplectic geometry are included. Examples of the Fedosov formalism in quantum mechanics are given.

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1 Introduction

The standard formulation of quantum mechanics in terms of a complex Hilbert space and linear operators is mainly applied for systems, whose classical limit may be described on a phase space $\mathbb{R}^{2n}$. There are two reasons for this situation: the fundamental operators $\hat{q}_i, \hat{p}_i, i = 1, \ldots, 2n$, are well defined by the Dirac quantization scheme [1] only in case $\mathbb{R}^{2n}$, and the operator orderings are based on the Fourier transform [2–4].

Unless we agree to weaken the foregoing assumptions [5] to deal with other quantum systems it is necessary to use geometric quantization [6, 7] or deformation quantization.
Classical mechanics is a physical theory which works perfectly on arbitrary differential manifolds. From this reason shortly after presenting a standard version of quantum theory, researchers began to look for an equivalent formulation of quantum mechanics based on differential geometry. The first complete version of quantum theory in the language of the theory of manifolds appeared in the middle of the XXth century, when Moyal [8] using previous work by Weyl [9], Wigner [10] and Groenewold [11] presented quantum mechanics as a statistical theory. His results are only valid for the case \( \mathbb{R}^{2n} \). However, the paper by Moyal contains the seminal ideas about deformation quantization, since the main result of this work is the substitution of the point-wise product of functions in phase space for a new product which is a formal power series in \( \hbar \) that we will precise later.

A modern version of Moyal’s deformation quantization on an arbitrary differential manifold was proposed by Bayen et al [12] in 1978. The mathematical structure of this formulation of the quantum theory is based, like Hamiltonian classical mechanics, on differential geometry of symplectic spaces. Observables are smooth real functions on a phase space and states are represented by functionals. Macro world appears in this formalism as the limit of the quantum reality for the Planck constant \( \hbar \) tending to \( 0^+ \).

The list of axioms constituting deformation quantization looks as follows:

i. a state of a physical system is described on a \( 2n \)-dimensional phase space \( M \),

ii. an observable is a real smooth function on \( M \),

iii. for every complex-valued smooth functions \( (f, g, h) \) of \( C^\infty(M) \) the \( \ast \)-product fulfills the following conditions:

(a) 
\[
f \ast g = \sum_{t=0}^{\infty} \left( \frac{i\hbar}{2} \right)^t M_t(f, g),
\]

where \( M_t(\cdot, \cdot) \) is a bidifferential operator on \( M \) (see definition later, formula (1.1)),

(b) at the classical limit \( \hbar \to 0^+ \)
\[
M_0(f, g) = f \cdot g;
\]

(c) the quasi-Dirac quantization postulate holds
\[
M_1(f, g) - M_1(g, f) = \{f, g\}_P,
\]

where \( \{\cdot, \cdot\}_P \) stands up for Poisson brackets;

(d) associativity also holds
\[
\sum_{t+u=s} (M_t(M_u(f, g), h) - M_t(f, M_u(g, h))) = 0, \quad \forall s \geq 0;
\]

(e) for the constant function equal to 1, we have
\[
M_t(1, f) = M_t(f, 1) = 0, \quad \forall t \geq 1.
\]
Let us comment the postulates written above. Assumptions (i) and (ii) say that the mathematical structure of deformation quantization is modeled after the classical Hamiltonian mechanics. Differences between classical and quantum descriptions appear at the level of product of functions and representation of states. In the quantum case the product must be noncommutative. The relation between this $\ast$-product and the usual point-wise multiplication of functions is established in (iii). The classical Lie algebra of functions is determined by the term $M_1$ from (iia). Associativity (iid) is analogous to the associativity in an algebra of linear operators in traditional quantum mechanics. Linearity of the bidifferential operators $M_t$ plus postulate (iie) expresses the fact that the measurement of a constant quantity does not interact with any other measurement.

Let us remember what is a bidifferential operator. Let $M$ be a $n$-dimensional differential manifold and $(U, \varphi)$ a local chart on it with coordinates $q = (q^1, \ldots, q^n)$. By a bidifferential operator $M_t(\cdot, \cdot) : \mathcal{C}^\infty(M) \times \mathcal{C}^\infty(M) \to \mathbb{C}$ we understand the map which locally in a chart $(U, \varphi)$ is a linear combination of terms like

$$l(q) \frac{\partial^m f(q)}{(\partial q^1)^{\alpha_1} \ldots (\partial q^n)^{\alpha_n}} \cdot \frac{\partial^r g(q)}{(\partial q^1)^{\beta_1} \ldots (\partial q^n)^{\beta_n}}, \quad l, f, g \in \mathcal{C}^\infty(M),$$

where $\alpha_1 + \cdots + \alpha_n = m, \beta_1 + \cdots + \beta_n = r$. The subindex ‘$t$’ in $M_t(\cdot, \cdot)$ denotes the highest derivative in the sum of elements of the kind (1.1) in $M_t(\cdot, \cdot)$.

Note that the assumptions of deformation quantization do not give a construction method of the $\ast$-product. In the simplest case when the phase space is $\mathbb{R}^{2n}$, there exists a correspondence between a wide class of linear operators acting on vectors from a Hilbert space $\mathcal{H}$ of the quantum system and functions on the phase space of it. More information about the class of operators and functions, for which such a correspondence exists, can be found in [13]. This relation, known as the Weyl application, in the 2-dimensional case takes the form [14]

$$\hat{F} = W(F) = \frac{1}{(2\pi)^2} \int_{\mathbb{R}^2} \alpha(h\lambda\mu) \exp (i[\lambda \hat{p} + \mu \hat{q}]) \hat{F}(\lambda, \mu) d\lambda d\mu \quad (1.2)$$

where $\hat{F}$ is a linear operator acting in $\mathcal{H}$, $\alpha(h\lambda\mu)$ is a function characterizing ordering of operators (however, in this work we will consider the case of Weyl ordering for which $\alpha(h\lambda\mu) = 1$), $p$ and $q$ are canonically conjugate coordinates on the phase space $\mathbb{R}^2$, $\hat{p}$ and $\hat{q}$ are self-adjoint operators representing momentum and position, respectively, and fulfilling the commutation relation

$$[\hat{q}, \hat{p}] = i\hbar$$

and, finally, $\hat{F}(\lambda, \mu)$ is the Fourier transform of the function $F(p,q)$ defined by

$$\hat{F}(\lambda, \mu) := \int_{\mathbb{R}^2} F(p,q) \exp (-i[\lambda p + \mu q]) dp dq.$$ 

The generalization of expression (1.2) to the case $\mathbb{R}^{2n}$ is straightforward.

Considering the $\ast$-product as the image of the product of operators by the inverse mapping $W^{-1}$ of the Weyl application $W$ (with $\alpha(h\lambda\mu) = 1$) we obtain that

$$F \ast G := W^{-1}(\hat{F} \cdot \hat{G}) = F \exp \left( \frac{i\hbar}{2} \mathcal{P} \right) G, \quad (1.3)$$
where $\mathbf{P}$ is the Poisson operator

$$
\mathbf{P} := \frac{\partial}{\partial q} \frac{\partial}{\partial p} - \frac{\partial}{\partial p} \frac{\partial}{\partial q} .
$$

The arrows indicate the acting direction of the partial differential derivatives. Explicitly

$$
F \ast G = \sum_{r=0}^{\infty} \sum_{l=0}^{r} \frac{1}{r!} \left( \frac{i\hbar}{2} \right)^r (-1)^l \frac{\partial^r F}{\partial q^{r-l} \partial p^l} \frac{\partial^l G}{\partial p^{r-l} \partial q^l} .
$$

The mapping $W^{-1}$ is called the Weyl correspondence and it constitutes an isomorphism between some algebra of linear operators and an algebra of functions on $\mathbb{R}^{2n}$. From this fact we deduce that the $\ast$-product defined as (1.3) is associative. Expressions (1.3) and (1.4) are true only in coordinate systems in which a symplectic form $\omega$ on $\mathbb{R}^2$ is canonical.

The construction presented above cannot be used in the general case because on an arbitrary manifold we are not able to introduce the Fourier transform, which is fundamental for defining the Weyl correspondence and the Weyl mapping.

As we mentioned before, Ref. [12] does not contain an explicit procedure for obtaining a $\ast$-product, only theorems of existence of $\ast$-products for symplectic manifolds are presented. However, other aspects of this paper have had a great influence in the modern development of deformation quantization. For instance, the deformation of the algebraic structure of the classical Poisson algebra $C^\infty(M)$ via the $\ast$-product provides an example of algebra deformation in the sense of Gerstenhaber [15]. Later, in the eighties, the algebraic deformation à la Gerstenhaber also appeared in relation to quantum groups [16]. Different aspects about deformation quantization and its applications can be found in [17]–[66]. For those who are interested in general aspects of quantization deformation we recommend the reviews [67–70].

The practical realization of deformation quantization on a symplectic manifold was proposed by Fedosov [71, 72]. His construction is purely geometric and it is based on the theory of fibre bundles [73]. Fedosov starts from some symplectic manifold endowed with a symmetric connection and lifts this connection to the so called Weyl algebra bundle. Next he builds some new flat connection in the Weyl bundle and shows how to operate on flat sections of the Weyl algebra bundle in order to define a $\ast$-product. In [74–78] one can find several considerations devoted to some aspects of the Fedosov formalism.

The aim of this paper is to present the Fedosov theory as an example of application of the theory of fibre bundles in physics. Thus, we analyze the construction of the $\ast$-product, paying attention to the geometrical aspects using the formalism of fibre bundles. All the final results presented here are known but we have obtained them applying consequently the fibre bundle methods and in many cases the proofs are new and easier than the original ones. This fact adds a pedagogical value to this work which can complement the monographies devoted to the Fedosov theory. As we mention before, we are mainly interested in the geometrical aspects of the Fedosov construction, for that reason we omit some long and technical proofs not connected directly with geometrical nature of the topic under study and that, moreover, can be found in the literature.
The paper is organized as follows. In section 2 we present some results about symplectic geometry that we will use later. The next section is devoted to the main mathematical structures involved on the Fedosov method. In section 4, starting from the construction of the Weyl algebra we obtain the Weyl bundle. This algebra is equipped with a symplectic connection, which is studied in detail in section 5. Making the symplectic connection flat and defining a one-to-one mapping between the formal series $\Pi(h, C^\infty(M))$ and flat sections of the Weyl bundle, we construct a noncommutative associative $\ast$-product on $M$, this is the content of the section 4. Some examples are presented in section 6.

## 2 Symplectic geometry

This part is devoted to review some aspects of the geometry of symplectic spaces. We present a procedure for defining parallel transport for such spaces and analyze the similarities and differences between Riemannian and symplectic geometry. The reader interested in the mathematical aspects of symplectic geometry can find a systematic presentation of this topic in [79]. Physical applications of the symplectic geometry are analyzed in [80–83].

**Definition 2.1.** A symplectic manifold $(M, \omega)$ is a manifold $M$ equipped with a nondegenerate closed 2-form $\omega$. The form $\omega$ is called a symplectic form.

It can be proved (see, for instance, [79]) that the dimension of the manifold $(M, \omega)$ is always even. Hereafter, we denote it by $\dim M = 2n$.

Let $(U_i, \varphi_i)$ be some local chart on $(M, \omega)$. In this chart the symplectic form may be written as

$$\omega = \omega_{ij} dq^i \wedge dq^j, \quad i, j = 1, 2, \ldots, 2n. \quad (2.1)$$

**Theorem 2.1.** (Darboux theorem) Let $(M, \omega)$ be a symplectic manifold. In the neighborhood of each point $p \in M$ there exist local coordinates $(q^1, \ldots, q^{2n})$, called canonical or Darboux coordinates, such that the form $\omega$ may be written by means of these coordinates as

$$\omega = dq^{n+1} \wedge dq^1 + dq^{n+2} \wedge dq^2 + \ldots + dq^{2n} \wedge dq^n. \quad (2.2)$$

An atlas $\{(U_\varnothing, \varphi_\varnothing)\}_{\varnothing \in I}$ consisting of Darboux charts is called the Darboux atlas on $(M, \omega)$.

As the form $\omega$ is nondegenerate, it establishes an isomorphism $I$ between tangent $T_pM$ and cotangent $T^*_pM$ spaces at each point $p \in M$. This isomorphism $I : T_pM \to T^*_pM$ is defined as follows.

**Definition 2.2.** The 1-form $I_X \in T^*_pM$, with $X \in T_pM$, satisfies the condition

$$I_X(Y) = \omega(Y, X), \quad \forall Y \in T_pM. \quad (2.3)$$

In a local chart $(U_i, \varphi_i)$ on $(M, \omega)$ in natural bases $\{dq^1, \ldots, dq^{2n}\}, \{\frac{\partial}{\partial q^1}, \ldots, \frac{\partial}{\partial q^{2n}}\}$ of $T^*_pM$ and $T_pM$, respectively, we can write that

$$(I_X)_r = \omega_{ir} X^i. \quad (2.4)$$
The inverse mapping $I^{-1} : T_p^* M \rightarrow T_p M$ can be obtained by

$$X^r = \omega^{ri}(I_X)_i,$$

where $\omega^{ri}$ is a covariant tensor for which the following relation holds

$$\omega^{ij}\omega_{jk} = \delta^i_k.$$

The Poisson bracket of two functions, $f, g \in C^\infty(M)$, is defined as

$$\{f, g\} = -\omega(I^{-1}(df), I^{-1}(dg)).$$

In local coordinates

$$\{f, g\} = \omega^{ij} \frac{\partial f}{\partial q^i} \frac{\partial g}{\partial q^j}.$$

From the closeness of the symplectic form $\omega$, i.e. $d\omega = 0$, we obtain the Jacobi identity

$$\{\{f, g\}, h\} + \{\{h, f\}, g\} + \{\{g, h\}, f\} = 0.$$

In fact, it is possible to define a Poisson structure on a manifold without introducing a symplectic form. In such cases by a Poisson structure on a manifold $M$ we understand an antisymmetric bilinear mapping $\{\cdot, \cdot\} : C^\infty(M) \times C^\infty(M) \rightarrow C^\infty(M)$ fulfilling the Jacobi identity (2.9) and the Leibniz rule

$$\{f, gh\} = g\{f, h\} + \{f, g\}h.$$

Poisson manifolds, i.e. pairs $(M, \{\cdot, \cdot\})$, are natural generalizations of symplectic manifolds [81].

The deformation quantization programme works on Poisson manifolds [53, 54, 62, 77], but we are going to continue our considerations for the case when the phase space of a system is some symplectic manifold.

There are two ways of building a phase space for a physical system:

- Starting from the configuration space of the system $\mathcal{M}$. At every point $p \in \mathcal{M}$ we can assign the cotangent space $T_p^* \mathcal{M}$. Taking the union $\bigcup_{p \in \mathcal{M}} T_p^* \mathcal{M}$ we obtain the cotangent bundle denoted by $T^* \mathcal{M}$. Such a bundle is equipped with the natural symplectic structure (for details see [83]). We define the phase space of the system as the cotangent bundle $T^* \mathcal{M}$. This algorithm is useful for systems without nonholonomic constraints and it is widely used in classical mechanics.

- Having a Lie group $G$, the symmetry group of the physical system, we construct orbits of this Lie group in the space of $K$—representation. These orbits can be interpreted as phase spaces for which the given Lie group is the group of symmetries [84].

We will follow the first procedure to construct a phase space. Let $(M, \omega)$ be a symplectic manifold. At every point $p \in M$ the tangent space $T_p M$ is assigned. Taking the union $\bigcup_{p \in M} T_p M$ we obtain the tangent bundle denoted by $TM$. 
**Definition 2.3.** A torsion-free connection $\nabla$ in the tangent bundle $TM$ is called symplectic if the induced connection $\nabla$ in $T^*M \otimes T^*M$ fulfills the condition

$$\nabla \omega = 0.$$

(2.10)

In a local chart $(U_i, \varphi_i)$ on $M$ for $\dim M = 2n$ the requirement (2.10) is equivalent to the system of $2n^2(2n-1)$ independent equations

$$\omega_{ij,k} := \frac{\partial \omega_{ij}}{\partial q^k} - \Gamma^l_{ik} \omega_{lj} - \Gamma^l_{jk} \omega_{il} = 0,$$

(2.11)

where $\Gamma^l_{ik}$ are coefficients of the connection $\nabla$. They are symmetric in lower indices because the connection is torsion-free. In a Darboux chart all the partial derivatives $\frac{\partial \omega_{ij}}{\partial q^k}$ vanish, so the equation system (2.11) is equivalent to the following one

$$-\Gamma_{jik} + \Gamma_{ijk} = 0$$

(2.12)

where

$$\Gamma_{ijk} := \omega_{il} \Gamma^l_{jk}.$$

**Corollary 2.1.** A connection $\nabla$ on the symplectic manifold $(M, \omega)$ is symplectic iff in every Darboux chart the coefficients $\Gamma_{ijk}$ are symmetric in all the indices $\{i, j, k\}$.

Unlike the Riemannian geometry on a symplectic manifold $(M, \omega)$ we can define many symplectic connections. The following statement holds.

**Theorem 2.2.** [72] The symplectic connection on a symplectic manifold $(M, \omega)$ is not unique. The set of coefficients

$$\tilde{\Gamma}_{ijk} := \Delta_{ijk} + \Gamma_{ijk}, \quad 1 \leq i, j, k \leq \dim M,$$

(2.13)

where $\Delta_{ijk}$ denotes the coefficients of a tensor symmetric with respect to indices $\{i, j, k\}$ and $\Gamma_{ijk}$ the coefficients of a symplectic connection on $(M, \omega)$, determines a symplectic connection on $(M, \omega)$.

**Proof.** Let us consider the change of coordinates $(q^1, \ldots q^{2n}) \equiv q \rightarrow (Q^1, \ldots Q^{2n}) \equiv Q$. The transformation rule for the coefficients $\Gamma_{ijk}$ is the following one

$$\Gamma_{ijk}(Q) = \Gamma_{rst}(q) \frac{\partial q^r}{\partial Q^i} \frac{\partial q^s}{\partial Q^j} \frac{\partial q^t}{\partial Q^k} + \omega_{rst}(q) \frac{\partial q^r}{\partial Q^i} \frac{\partial q^s}{\partial Q^j} \frac{\partial^2 q^t}{\partial Q^k \partial Q^r \partial Q^s}.$$

(2.14)

For the tensor $\Delta_{ijk}$ the relation holds

$$\Delta_{ijk}(Q) = \Delta_{rst}(q) \frac{\partial q^r}{\partial Q^i} \frac{\partial q^s}{\partial Q^j} \frac{\partial q^t}{\partial Q^k}.$$

(2.15)

From (2.14) and (2.13) we can see that under the change of coordinates $q \rightarrow Q$ the sum $\Gamma_{ijk} + \Delta_{ijk}$ transforms under the rule (2.14), so $\tilde{\Gamma}_{ijk}$ is a connection.
In a Darboux chart, if \( \Gamma_{ijk} \) and \( \Delta_{ijk} \) are symmetric with respect to indices \( \{i,j,k\} \), the same happens also for the sum \( \tilde{\Gamma}_{ijk} = \Gamma_{ijk} + \Delta_{ijk} \). It means that \( \tilde{\Gamma}_{ijk} \) is not only a connection but also a symplectic connection.

Of course, the tensor \( \Delta_{ijk} \) is symmetric in any coordinate system.

Note that locally symplectic connections exist on any symplectic manifold. The construction of a symplectic connection on the whole manifold \( (M,\omega) \) can be done using a \( C^\infty \)-partition of unity.

**Definition 2.4.** [85] A manifold \( M \) admits \( C^\infty \)-partitions of unity if, given a locally finite open cover \( \{U_i\}_{i \in I} \), there exists a family of \( C^\infty \)-maps \( \psi_i : M \to [0,1] \) such that the set of points at which \( \psi_i \) does not vanish, \( \text{supp}(\psi_i) \), is contained in \( U_i \), i.e. \( \text{supp}(\psi_i) \subset U_i \), and \( \sum_{i \in I} \psi(p) = 1 \) for all \( p \in M \).

**Theorem 2.3.** [85] Every manifold \( M \) admits \( C^\infty \)-partitions of unity.

Since theorem 2.3 holds, on an arbitrary symplectic manifold \( (M,\omega) \) we build the symplectic connection on \( (M,\omega) \) following the algorithm:

1. We cover the manifold \( (M,\omega) \) with a Darboux atlas \( \{(U_i,\varphi_i)\}_{i \in I} \). According to theorem 2.3 there exists a \( C^\infty \)-partition of unity \( \{\psi_i\}_{i \in I} \) compatible with \( \{U_i\}_{i \in I} \).
2. In each chart \( (U_i,\varphi_i) \) we define a symplectic connection by giving coefficients \( (\Gamma_{jkl})_i \) symmetric in all the indices.
3. The coefficients of the symplectic connection \( \nabla \) on \( (M,\omega) \) at an arbitrary fixed point \( p \in M \) are expressed by

\[
\Gamma_{jkl}(p) := \sum_{i \in I} \psi_i(p)(\Gamma_{jkl})_i.
\]

**Definition 2.5.** Let \( (M,\omega) \) be a symplectic manifold with a symplectic connection \( \nabla \). The triad \( (M,\omega,\nabla) \) is called Fedosov manifold.

The curvature of a symplectic connection \( \nabla \) is characterized by a curvature tensor.

**Definition 2.6.** The curvature tensor \( R \) of a symplectic connection \( \nabla \) in the tangent bundle \( TM \) is a mapping \( R : TM \times TM \times TM \to TM \) fulfilling the relation

\[
R(X,Y)Z = \nabla_X \nabla_Y Z - \nabla_Y \nabla_X Z - \nabla_{[X,Y]}Z, \quad \forall X,Y,Z \in TM.
\] (2.16)

By \( [X,Y] \) we denote, as usually, the Lie bracket of the vector fields \( X \) and \( Y \). In the natural basis \( \{\frac{\partial}{\partial q^i}, \ldots, \frac{\partial}{\partial q^n}\} \) of \( TM \) the components of the curvature tensor are expressed as

\[
R \left( \frac{\partial}{\partial q^l}, \frac{\partial}{\partial q^m} \right) \frac{\partial}{\partial q^i} = R^{m}_{ijkl} \frac{\partial}{\partial q^m}.
\] (2.17)

In terms of the Christoffel symbols

\[
R^l_{ijk} = \frac{\partial \Gamma^l_{ki}}{\partial q^j} - \frac{\partial \Gamma^l_{ij}}{\partial q^k} + \Gamma^m_{ki} \Gamma^l_{jm} - \Gamma^m_{ij} \Gamma^l_{km}.
\] (2.18)
Note that for various symplectic connections we may have different curvature tensors. For instance, the same symplectic manifold \((M, \omega)\) can be equipped with flat (all the components \(R_{l ijk}\) of curvature tensor vanish) or nonvanishing curvature.

**Theorem 2.4.** Let us consider two Fedosov manifolds \((M, \omega, \nabla)\) and \((M, \omega, \tilde{\nabla})\) for which the relation holds \(\tilde{\nabla} = \nabla + \Delta\) (see theorem 2.2). The curvatures \(R\) and \(\tilde{R}\) fulfill the equation

\[
\tilde{R} = R + r(\Delta),
\]

where \(r(\Delta)\) is some tensor depending only on \(\Delta\).

**Proof.** Since locally in a chart \((U_i, \varphi_i)\) the equality \(\tilde{\Gamma}^i_{jk} = \Gamma^i_{jk} + \omega^s\Delta_{skj}\) is fulfilled, from the relation (2.18) we obtain that

\[
\tilde{R}^l_{ijk} = \frac{\partial \Gamma^l_{ki}}{\partial q^j} + \frac{\partial \omega^s\Delta_{skj}}{\partial q^j} - \frac{\partial \omega^s\Delta_{sij}}{\partial q^k} + \omega^m\Delta_{skl}\omega^z\Delta_{zjm} - \omega^m\Delta_{sjl}\omega^z\Delta_{zkm}.
\]

It is always possible to choose the chart \((U_i, \varphi_i)\) in such a way that at an arbitrary fixed point \(p \in M\) all the Christoffel symbols \(\Gamma^l_{jm}\) vanish. It means that at the point \(p\) we have

\[
\tilde{R}^l_{ijk} = R^l_{ijk} + r(\Delta)^l_{ijk},
\]

where

\[
r(\Delta)^l_{ijk} = \frac{\partial \omega^s\Delta_{skj}}{\partial q^j} - \frac{\partial \omega^s\Delta_{sij}}{\partial q^k} + \omega^m\Delta_{skl}\omega^z\Delta_{zjm} - \omega^m\Delta_{sjl}\omega^z\Delta_{zkm}.
\]

Note that although \(r(\Delta)^l_{ijk}\) looks like a curvature tensor, in fact, it is not a curvature because \(\Delta\) is a tensor not a connection.

Lowering the upper index in \(R^l_{ijk}\) we define a new tensor

\[
R_{ijkl} := \omega_{mi}R^m_{ijkl}.
\]

This tensor has the following properties [86, 87]:

1. The curvature tensor is antisymmetric in the last two indices

\[
R_{ijkl} = -R_{ijkl}.
\]

Indeed, from (2.18) and the fact that the connection \(\nabla\) is symmetric we can see that \(R^l_{ijk} = -R^l_{ikj}\). Making the contraction (2.22) we obtain (2.23).
2.- The first Bianchi identity holds

\[ R_{ijkl} + R_{iklj} + R_{iljk} = 0. \]  \hspace{1cm} (2.24)

Effectively, the identity (2.24) is equivalent to the relation

\[ R^i_{jkl} + R^i_{kli} + R^i_{lij} = 0. \]  \hspace{1cm} (2.25)

It means that

\[ R \left( \frac{\partial}{\partial q^k}, \frac{\partial}{\partial q^l} \right) \frac{\partial}{\partial q^j} + R \left( \frac{\partial}{\partial q^l}, \frac{\partial}{\partial q^k} \right) \frac{\partial}{\partial q^i} + R \left( \frac{\partial}{\partial q^k}, \frac{\partial}{\partial q^i} \right) \frac{\partial}{\partial q^l} = 0. \]  \hspace{1cm} \text{(2.28)}

On the other hand, the Lie brackets

\[ \left[ \frac{\partial}{\partial q^j}, \frac{\partial}{\partial q^k} \right] = 0, \quad j, k = 1, \ldots, 2n. \]  \hspace{1cm} (2.26)

Moreover, the connection \( \nabla \) is symmetric, so that

\[ \nabla \frac{\partial}{\partial q^k} \frac{\partial}{\partial q^j} = \nabla \frac{\partial}{\partial q^j} \frac{\partial}{\partial q^k}, \quad j, k = 1, \ldots, 2n. \]  \hspace{1cm} (2.27)

From the definition of the curvature (2.16), involving (2.26) we rewrite the l.h.s. of (2.25) in the form

\[ \nabla \frac{\partial}{\partial q^j} \nabla \frac{\partial}{\partial q^k} - \nabla \frac{\partial}{\partial q^k} \nabla \frac{\partial}{\partial q^j} + \nabla \frac{\partial}{\partial q^j} \nabla \frac{\partial}{\partial q^l} - \nabla \frac{\partial}{\partial q^l} \nabla \frac{\partial}{\partial q^j}. \]  \hspace{1cm} (2.28)

From (2.27) we can see that the foregoing expression vanishes.

3.- The second Bianchi identity is verified, i.e.

\[ R_{mjkl;i} + R_{mjk;l} + R_{mji;k} = 0. \]  \hspace{1cm} (2.29)

Indeed, this identity is equivalent to the relation

\[ R^m_{jkli} + R^m_{jkil} + R^m_{jilk} = 0. \]  \hspace{1cm} (2.30)

It is always possible to consider locally such a chart on \((M, \omega, \nabla)\) that at an arbitrary fixed point \(p \in M\) all the coefficients \(\Gamma^{ij}_{jk}\) disappear. If we compute all the covariant derivatives from (2.30) at \(p\), we obtain that the l.h.s. of (2.30) equals

\[ \frac{\partial^2 \Gamma^m_{ij}}{\partial q^i \partial q^j} - \frac{\partial^2 \Gamma^m_{kj}}{\partial q^k \partial q^j} - \frac{\partial^2 \Gamma^m_{ji}}{\partial q^j \partial q^i} + \frac{\partial^2 \Gamma^m_{ij}}{\partial q^q \partial q^i} - \frac{\partial^2 \Gamma^m_{ij}}{\partial q^q \partial q^i} = 0. \]  \hspace{1cm} \text{(2.31)}

It is easy to see that in fact properties 1–3 hold for any curvature tensor of a symmetric connection. Now we are going to present some features of the symplectic curvature tensors exclusively \[86\].
4.- The symplectic curvature tensor $R_{ijkl}$ is symmetric in the first two indices, i.e.

$$R_{ijkl} = R_{jikl}. \quad (2.31)$$

The proof is as follows. From the Darboux theorem we can always find a system of coordinates in which the symplectic tensor $\omega_{ij}$ is locally constant. In such a chart we have

$$R_{ijkl} = \frac{\partial \Gamma_{ilj}}{\partial q^k} - \frac{\partial \Gamma_{ijk}}{\partial q^l} + \omega^{zu}_{\Gamma_{ulj}} \Gamma_{ikz} - \omega^{zu}_{\Gamma_{uik}} \Gamma_{jlz}. \quad (2.32)$$

$$R_{jikl} = \frac{\partial \Gamma_{jli}}{\partial q^k} - \frac{\partial \Gamma_{jik}}{\partial q^l} + \omega^{uz}_{\Gamma_{uli}} \Gamma_{jkz} - \omega^{uz}_{\Gamma_{uik}} \Gamma_{jlz}. \quad (2.33)$$

Permuting repeated indices $z \leftrightarrow u$ in (2.33) we see that

$$R_{jikl} = \frac{\partial \Gamma_{jli}}{\partial q^k} - \frac{\partial \Gamma_{jik}}{\partial q^l} + \omega^{uz}_{\Gamma_{uli}} \Gamma_{jkz} - \omega^{uz}_{\Gamma_{uik}} \Gamma_{jlz}.$$

Remembering that $\omega^{zu} = -\omega^{uz}$ and using the fact that in Darboux coordinates the coefficients $\Gamma_{jli}$ are symmetric in all the indices we obtain that $R_{ijkl} = R_{jikl}$. ■

5.- For a symplectic curvature tensor the following relation holds

$$R_{ijkl} + R_{jkli} + R_{klij} + R_{lijk} = 0. \quad (2.34)$$

Indeed, from the first Bianchi identity (2.24) we get

$$R_{ijkl} + R_{iljk} + R_{iklj} = 0,$$
$$R_{jikl} + R_{jikl} + R_{jkl} = 0,$$
$$R_{klij} + R_{klij} + R_{kli} = 0,$$
$$R_{lijk} + R_{lijk} + R_{lij} = 0. \quad (2.35)$$

From (2.23) and (2.31) adding formulas (2.35) we conclude that the equation (2.32) is always true for any symplectic curvature. ■

Let us compute the number of independent components of the tensor $R_{ijkl}$. As usually, we assume that $\dim M = 2n$. From the properties of antisymmetry (2.23) and symmetry (2.31) we obtain that the tensor $R_{ijkl}$ has $n^2(2n + 1)(2n - 1)$ independent elements. The first Bianchi identity (2.24) is a new constraint if and only if all the indices $j, k, l$ are different in $R_{ijkl}$. It means that (2.24) provides us $2n \left( \begin{array}{c} 2n \\ 3 \end{array} \right)$ new equations. The property (2.31) can no be reduced by the symmetry of the tensor or by the first Bianchi identity only if all the indices $i, j, k, l$ are different, so we have $\left( \begin{array}{c} 2n \\ 4 \end{array} \right)$ independent conditions. Hence, the symplectic curvature tensor $R_{ijkl}$ has

$$n(2n - 1) \left( n(2n + 1) - \frac{2}{3} n(2n - 2) - \frac{1}{12} (2n - 2)(2n - 3) \right).$$
independent components. For example, on a 2-dimensional Fedosov manifold \((M, \omega, \nabla)\) its curvature tensor has 3 independent components. For \(\dim M = 4\) the amount of independent elements of \(R_{ijkl}\) increases to 47. When \(M\) is a Riemannian manifold, the number of independent components of a curvature tensor would be equal to 1 and 20, respectively.

Apart from the curvature tensor \(R_{ijkl}\) the geometry of a symplectic space can be characterized by a Ricci tensor and a scalar of curvature.

**Definition 2.7.** The Ricci tensor on a Fedosov manifold \((M, \omega, \nabla)\) is defined by

\[
K_{ij} := \omega^{kl} R_{likj} = R_{iklj}. \tag{2.36}
\]

The Ricci tensor is symmetric, i.e.

\[
K_{ij} = K_{ji}. \tag{2.37}
\]

Indeed, from \((2.23)\), \((2.31)\) and \((2.34)\) we can write

\[
\omega^{ki}(R_{ijkl} + R_{jkl} + R_{kli} + R_{lijk}) = K_{jl} + K_{jl} - K_{lj} - K_{lj} = 0.
\]

So, the equality \((2.37)\) holds.

**Corollary 2.2.** On any Fedosov manifold

\[
\omega^{ji} K_{ij} = 0. \tag{2.38}
\]

This conclusion is a straightforward consequence of the fact that the Ricci tensor \(K_{ij}\) is symmetric and \(\omega^{ji}\) antisymmetric. From \((2.38)\) we can see that the scalar of any symplectic curvature defined as \(K := \omega^{ji} K_{ij}\) is always 0.

## 3 The Weyl bundle

In this section we present the construction of a bundle, which plays a fundamental role in the definition of the \(\ast\)-product on a symplectic manifold. Let us start defining a formal series over a vector space.

**Definition 3.8.** \([75]\) Let \(\lambda\) be a fixed real number and \(V\) some vector space. A formal series in the formal parameter \(\lambda\) is each expression of the form

\[
v[[\lambda]] = \sum_{i=0}^{\infty} \lambda^i v_i, \quad \forall v_i \in V. \tag{3.1}
\]

The set of formal series \(v[[\lambda]]\) constitutes a vector space. Addition means vector summation of the elements of the same power of \(\lambda\), and multiplication by a scalar, \(a\), is multiplication of each vector standing on the r.h.s. of \((3.1)\) by \(a\), i.e.

\[
u[[\lambda]] + v[[\lambda]] = \sum_{i=0}^{\infty} \lambda^i (u_i + v_i)
\]

\[
a \cdot v[[\lambda]] = \sum_{i=0}^{\infty} \lambda^i (av_i).
\]
The vector space of formal series over the vector space $V$ in the parameter $\lambda$ will be denoted by $V[[\lambda]]$ and can be considered as a direct sum

$$V[[\lambda]] = \bigoplus_{i=1}^{\infty} V_i, \quad V_i = V \quad \forall i. \quad (3.2)$$

Let $(M, \omega)$ be a symplectic manifold and $T^*_p M$ the cotangent space to $M$ at the point $p$ of $M$. The space $(T^*_p M)^l$ is a symmetrized tensor product of $T^*_p M \otimes^{l-\text{times}} \ldots \otimes T^*_p M$. It is spanned by

$$v_{i_1} \otimes \ldots \otimes v_{i_l} := \frac{1}{l!} \sum_{\text{all permutations}} v_{\sigma_{i_1}} \otimes \ldots \otimes v_{\sigma_{i_l}}, \quad (3.3)$$

where $v_{i_1}, \ldots, v_{i_l} \in T^*_p M$.

**Definition 3.9.** A preweyl vector space $P_p^* M$ at the point $p \in M$ is the direct sum

$$P_p^* M := \bigoplus_{l=0}^{\infty} (T^*_p M)^l. \quad \text{(3.4)}$$

We introduce the formal series over the preweyl vector space as follows.

**Definition 3.10.** A Weyl vector space $P_p^* M[[\hbar]]$ is the vector space of formal series over the preweyl vector space $P_p^* M$ in the formal parameter $\hbar$.

For further physical applications we usually identify $\hbar$ with the Planck constant. The elements of $P_p^* M[[\hbar]]$ can be written in the form

$$v_p[[\hbar]] = \sum_{k=0}^{\infty} \sum_{l=0}^{\infty} \hbar^k a_{k,i_1 \ldots i_l}, \quad a_{k,i_1 \ldots i_l} \in P_p^* M. \quad (3.4)$$

For $l = 0$ we have just the sum $\sum_{k=0}^{\infty} \hbar^k a_k$.

**Definition 3.11.** [72] The degree of the component $a_{k,i_1 \ldots i_l}$, $\deg a_{k,i_1 \ldots i_l}$ of the Weyl vector space $P_p^* M[[\hbar]]$ is $2k + l$.

Our aim is to define a product (denoted by $\circ$) of elements of the Weyl space $P_p^* M[[\hbar]]$ which equips the Weyl space with an algebra structure. Such a product must give a symmetric tensor. Moreover, we require that $\deg(a \circ b) = \deg a + \deg b$.

Let us assume that the elements of the Weyl vector space are written in a natural basis constructed in terms of symmetric tensor products of $\{dq^1, \ldots, dq^{2n}\}$, where $\dim M = 2n$, and $X_p \in T_p M$ is some fixed vector of the space $T_p M$ tangent to $M$ at the point $p$. Let us denote the components of $X$ in the basis $\{\frac{\partial}{\partial q^1}, \ldots, \frac{\partial}{\partial q^{2n}}\}$ by $X^i_p$.

It is obvious that for every $a_{k,i_1 \ldots i_l} \in P_p^* M[[\hbar]]$

$$a_{k,i_1 \ldots i_l}(X_p, \ldots, X_p) = a_{k,i_1 \ldots i_l} X^i_p \ldots X^i_p$$

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is a complex number and we can handle $a_{k,i_1\ldots i_l}(X_p, \ldots, X_p)$ as a polynomial of $l$-th degree in the components of the vector $X_p$. We extend this observation to each element of the Weyl algebra $P^*_pM[[\hbar]]$ and consider $v_p[[\hbar]](X_p)$ like a function of $X_p^1, \ldots, X_p^{2n}$ of the form

$$v_p[[\hbar]](X_p) = \sum_{k=0}^{\infty} \sum_{l=0}^{\infty} \hbar^k a_{k,i_1\ldots i_l} X_p^{i_1} \cdots X_p^{i_l}. \quad (3.5)$$

We do not define any topology in the space of functions of the kind (3.5).

By the derivative $\frac{\partial v_p[[\hbar]]}{\partial X_{i_p}}$ we understand the derivative of the sum (3.5) as function of $X_p^1, \ldots, X_p^{2n}$.

**Definition 3.12.** [71, 72] The product $\circ : P^*_pM[[\hbar]] \times P^*_pM[[\hbar]] \to P^*_pM[[\hbar]]$ of two elements $a, b \in P^*_pM[[\hbar]]$ is an element $c \in P^*_pM[[\hbar]]$ such that for each $X_p \in T_pM$ the following equality holds

$$c(X_p) = a(X_p) \circ b(X_p) := \sum_{t=0}^{\infty} \frac{1}{t!} \left( \frac{i\hbar}{2} \right)^t \omega^{i_1j_1} \cdots \omega^{i_lj_l} \partial^t a(X_p) \partial X_p^{i_1} \cdots \partial X_p^{i_l} \partial^t b(X_p) \partial X_p^{j_1} \cdots \partial X_p^{j_l}. \quad (3.6)$$

The pair $(P^*_pM[[\hbar]], \circ)$ is a noncommutative associative algebra called the Weyl algebra and denoted by $P^*M_p[[\hbar]]$. Let us simply enumerate some properties of the $\circ$-product:

1. The $\circ$-product is independent of the chart. This result is a straightforward consequence of the fact that all the elements appearing on the r.h.s. of the formula (3.6) are scalars.

2. The $\circ$-product is associative but in general nonabelian. The associativity of the $\circ$-product may be roughly explained in the following way. The elements $a(X_p)$ and $b(X_p)$ are polynomials in $X^i$'s. Let us formally substitute $X^i \to q^i$, where by $q^i$ we denote the cartesian coordinates on the symplectic space ($\mathbb{R}^{2n}, \omega$). After this substitution the product for polynomials in variables $q^i$'s defined like (3.6) is a associative Moyal product (see the relation (1.3)). Indeed, in Darboux coordinates in the 2-dimensional case with the symplectic form $\omega = dq^2 \land dq^1$, the formula (3.6) may be written as

$$a \circ b = \sum_{t=0}^{\infty} \frac{1}{t!} \left( \frac{i\hbar}{2} \right)^t \sum_{l=0}^{t} (-1)^l \frac{\partial^t a(X_p)}{\partial (X_p^1)^{t-l} \partial (X_p^2)^l} \frac{\partial^t b(X_p)}{\partial (X_p^1)^l \partial (X_p^2)^{t-l}}$$

which is exactly (1.4).

3. The degree verifies that

$$\deg (a \circ b) = \deg a + \deg b. \quad (3.7)$$

This statement, according to the definition (3.6), is obvious.
In order to enlighten the above construction of the $\circ$-product let us consider the following example. Let $(M, \omega)$ be a 2-dimensional symplectic manifold. From the Darboux theorem [24] we can choose a chart $(U_i, \varphi_i)$ in such a manner that at a fixed point $p \in M$ and in some neighborhood of it we have that
\[ \omega = dq^2 \land dq^1. \]
It means that a chart $(U_i, \varphi_i)$ in a natural basis $\left\{ \frac{\partial}{\partial q^i}, \frac{\partial}{\partial q^j} \right\}$ the tensor $\omega^{ij}$ takes the form
\[ \omega^{ij} = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}. \]
Let us consider the $\circ$-product of elements $a = a_{0,11}$ and $b = b_{0,12}$, where the symbol $\left[ \ldots \right]$ denotes the symmetrization in the indices inside the bracket. From (3.6) we obtain
\[
a(X) \circ b(X) = a_{0,11}X^1X^1 \circ b_{0,12}X^1X^2 = a_{0,11}b_{0,12}X^1X^1X^2 + \frac{1}{\hbar} \left( \text{ii} \right) \cdot 1 \cdot 2a_{0,11}b_{0,12}X^1X^1. \]
So, finally
\[ c_{0,[112]} = a_{0,11}b_{0,12} \quad c_{1,[11]} = i \cdot a_{0,11}b_{0,12}. \]

We have worked until now at an arbitrary fixed point $p$ of the symplectic manifold $(M, \omega)$. But our aim is to define the $*$-product on the whole manifold. To achieve it we need to introduce a new object, the so called Weyl algebra bundle.

Let us start reminding the definition of a fibre bundle.

**Definition 3.13.** [88] A (differentiable) fibre bundle $E \equiv (E, \pi, M, F, G)$ consists of the following elements: (1) the differentiable manifolds $E, M$ and $F$ called the total space, the base space and the fibre (or typical fibre), respectively; (2) a surjection $\pi : E \to M$ called the projection. The inverse image $\pi^{-1}(p) \equiv F_p \cong F$ is called the fibre at $p \in M$; (3) a Lie group $G$ called the structure group, which acts on $F$ at the left; (4) a set of open coverings $\{U_i\}$ of $M$ with a diffeomorphism $\varphi_i : U_i \times F \to \pi^{-1}(U_i)$ such that $\pi \varphi_i(p, f) = p$. The map $\varphi_i$ is called the local trivialization since $\varphi_i^{-1}$ maps $\pi^{-1}(U_i)$ onto the direct product $U_i \times F$; (5) if we write $\varphi_i(p, f) = \varphi_{i,p}(f)$, the map $\varphi_{i,p} : F \to F_p$ is a diffeomorphism. On $U_i \cap U_j \neq \emptyset$, we require that $t_{ij}(p) := \varphi_{i,p}^{-1} \varphi_{j,p} : F \to F$ be an element of $G$. Then $\varphi_i$ and $\varphi_j$ are related by a smooth map $t_{ij} : U_i \cap U_j \to G$ as $\varphi_j(p, f) = \varphi_i(p, t_{ij}(p)f)$.

The elements $\{t_{ij}\}$ are called the transition functions.

A section of the fibre bundle is a map $s : M \to E$ such that $(\Pi \circ s)(p) = p$ for every $p$ of $M$. The space of all the (smooth) sections will be denoted by $C^\infty(M, E)$ or simply $C^\infty(E)$.

A differential fibre bundle where the typical fibre $F$ and the fibres $F_p$ are vector spaces is called a vector bundle. Obviously, if $F$ and the fibres $F_p$ have the structure of algebra we have an algebra bundle.
Definition 3.14. The collection of all the Weyl algebras, i.e.

\[ \mathcal{P}^*\mathcal{M}[[\hbar]] = \bigcup_{p \in M} \mathcal{P}^*\mathcal{M}_p[[\hbar]], \]  

is a vector bundle which is also an algebra bundle and it is called the Weyl algebra bundle.

The structure of the Weyl algebra bundle looks as follows:

- The set \( \mathcal{P}^*\mathcal{M}[[\hbar]] \) is the total space.
- The symplectic space \((M, \omega)\) is the base space.
- The fibre is the Weyl algebra \( \mathcal{P}^*\mathcal{M}_p[[\hbar]] \). We do not introduce a new symbol but now \( \mathcal{P}^*\mathcal{M}_p[[\hbar]] \) is not connected with any point. Moreover, \( \mathcal{P}^*\mathcal{M}_p[[\hbar]] \) is a Weyl vector space and \( \omega_{ij}, \ i,j = 1, \ldots, 2n \), are the coefficients of a fixed nondegenerate antisymmetric tensor.
- The projection \( \pi : \mathcal{P}^*\mathcal{M}[[\hbar]] \to M \) is defined by \( \Pi(v_p[[\hbar]]) = p \).
- \( GL(2n, \mathbb{R}) \) is the group of real automorphisms of \( T^*_pM \). The structure group of the fibre is

\[ G := \bigoplus_{l=0}^{\infty} (GL(2n, \mathbb{R}) \otimes \ldots \otimes GL(2n, \mathbb{R})) \text{ } l\text{-times}. \]  

For \( l = 0 \) we have just the identity transformation. The tensor \( \omega_{ij} \) transforms under the group \( GL(2n, \mathbb{R}) \otimes GL(2n, \mathbb{R}) \). Moreover, if the tensor \( a_{ij} \) transforms under the element \( g \in GL(2n, \mathbb{R}) \otimes GL(2n, \mathbb{R}) \), then also \( \omega_{ij} \) transforms under \( g^{-1} \in GL(2n, \mathbb{R}) \otimes GL(2n, \mathbb{R}) \).

- Let \((U_i, \phi_i)\) be a chart on \( M \). The local trivialization \( \varphi_i \) is a diffeomorphism which assigns to every point of \( \mathcal{P}^*\mathcal{M}[[\hbar]] \) the point \( p \) of \( M \) and to the element of \( \mathcal{P}^*\mathcal{M}_p[[\hbar]] \) its coordinates in the natural basis \( \left\{ \frac{\partial}{\partial q_{i1}} \otimes \ldots \otimes \frac{\partial}{\partial q_{il}} \right\} \) determined by \( \phi_i \).

Other kind of fibre bundle used in this work is the cotangent bundle.

Definition 3.15. [88] The cotangent bundle \( T^*M \) of a differentiable manifold \( M \) is

\[ T^*M \equiv \bigcup_{p \in M} T^*_pM, \]  

where \( T^*_pM \) is the cotangent space at the point \( p \) of \( M \).

Let \( \Lambda^k \) be the vector bundle

\[ (T^*M \otimes \ldots \otimes T^*M, \pi, M, T^*M, GL(2n, \mathbb{R}) \otimes \ldots \otimes GL(2n, \mathbb{R})) \]  

\( k\)-times + antisymmetrization
of $k$-forms on $M$. Taking the direct sum of tensor products of bundles

$$\mathcal{P}^*\mathcal{M}^{[[\hbar]]}\Lambda = \bigoplus_{k=0}^{2n} \mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^k$$

we obtain a new algebra bundle. A product can be defined in it in terms of the $\circ$-product of elements of the Weyl algebra $\mathcal{P}^*\mathcal{M}_{p}^{[[\hbar]]}$ and the external product ($\wedge$) of forms. We will denote the new product also by ‘$\circ$’. In a local chart $(U_i, \varphi_i)$ the elements of $\mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^k$ are smooth tensor fields of the kind

$$a_{c,i_1...i_l,j_1...j_k}(q^1, \ldots, q^{2n})$$

These objects $a_{c,i_1...i_l,j_1...j_k}$ are symmetric in the indices $(i_1 \ldots i_l)$ (as elements of the Weyl algebra) and antisymmetric in $(j_1, \ldots, j_k)$ (as forms). For simplicity we will omit coordinates $(q^1, \ldots, q^{2n})$. The elements of $\mathcal{P}^*\mathcal{M}^{[[\hbar]]}\Lambda$ can be seen as forms with values in the Weyl algebra.

In the special case when $a \in \Lambda^0$ is a smooth function on $M$, we obtain

$$a \circ b = a \cdot b = a \wedge b, \quad \forall b \in \mathcal{P}^*\mathcal{M}^{[[\hbar]]}\Lambda.$$  

**Definition 3.16.** [78] The commutator of two forms $a \in \mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^{k_1}$ and $b \in \mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^{k_2}$ is a form belonging to $\mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^{k_1+k_2}$ such that

$$[a, b] := a \circ b - (-1)^{k_1 \cdot k_2} b \circ a. \quad (3.11)$$

We point out some properties of the commutator of forms with values in the Weyl algebra:

1. The straightforward consequence of (3.11) is the equality

$$[b, a] = (-1)^{k_1 \cdot k_2 + 1} [a, b]. \quad (3.12)$$

2. The Jacobi identity

$$(-1)^{k_2(k_1+k_3)} [a, [b, c]] + (-1)^{k_1(k_2+k_3)} [c, [a, b]] + (-1)^{k_3(k_1+k_2)} [b, [c, a]] = 0 \quad (3.13)$$

holds for every $a \in \mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^{k_1}$, $b \in \mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^{k_2}$ and $c \in \mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^{k_3}$. It is a simple matter of computation to prove this identity.

3.

$$[a, b \circ c] = (-1)^{k_1 k_2} b \circ [a, c] + [a, b] \circ c \quad (3.14)$$

for all $a \in \mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^{k_1}$, $b \in \mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^{k_2}$ and $c \in \mathcal{P}^*\mathcal{M}^{[[\hbar]]} \otimes \Lambda^{k_3}$.

**Proof.** Indeed,

$$[a, b \circ c] = a \circ b \circ c - (-1)^{k_1(k_2+k_3)} b \circ c \circ a. \quad (3.15)$$

Developing the r.h.s. of (3.14) we get

$$(-1)^{k_1 k_2} b \circ [a, c] = (-1)^{k_1 k_2} b \circ a \circ c - (-1)^{k_1 k_2} (-1)^{k_1 k_3} b \circ c \circ a \quad (3.16)$$

and

$$[a, b] \circ c = a \circ b \circ c - (-1)^{k_1 k_2} b \circ a \circ c. \quad (3.17)$$

Summing (3.16) and (3.17) we recover (3.14).
4. The commutator $[a,b]$ of two forms $a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^m$ and $b \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^n$ contains only terms with odd number of derivatives in $X^i$.

5. The commutator $[a,b]$ of two real forms $a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^m$ and $b \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^n$ is purely imaginary.

Proofs of these last two properties can be found in [89].

## 4 Connections in the bundle $\mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda$

Let $(M, \omega, \nabla)$ be a Fedosov manifold equipped with a symplectic connection $\nabla$, whose coefficients in the Darboux atlas $\{(U_\varphi, \varphi)\}_{\varphi \in \mathcal{I}}$ are $\Gamma_{i\ell k}$. The connection allows us to know how to transport geometrical quantities parallelly on $M$.

Let us define the covariant derivative of a tensor $a_{t_{i_1}i_2...j_1j_2}$ with respect the variable $q^s$. Following [88] we can write

$$
\partial_{s}a_{t_{i_1}i_2...j_1j_2} := \frac{\partial a_{t_{i_1}i_2...j_1j_2}}{\partial q^s} - \omega^{ur}\Gamma_{rsi_1t_{i_1}i_2...j_1j_2} - \ldots
$$

The covariant derivative $\partial_{s}a$ increases the covariance of the tensor $a_{t_{i_1}i_2...j_1j_2}$, so the object $\partial_{s}a_{t_{i_1}i_2...j_1j_2}$ is a tensor of the range $(0,l + k + 1)$. On the other hand, the covariant derivative $\partial_{s}a_{s_{i_1}...i_2j_1...j_2}$ is again symmetric in indices $(i_1, \ldots, i_l)$ and antisymmetric in $(j_1, \ldots, j_k)$. It means that the operation $\partial_{s}$ changes neither the degree of the Weyl algebra element nor the degree of the form. The new index ‘$s$’ does not generate inner symmetries in $\partial_{s}a_{t_{i_1}i_2...j_1j_2}$. The covariant derivative is non symmetric in $(s, i_1, \ldots, i_l)$ and non antisymmetric in $(s,j_1,\ldots,j_k)$. The straightforward consequence of this fact is that $\partial_{s}a_{t_{i_1}i_2...j_1j_2}$ is not an element of $\mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda$.

Such a result is not compatible with our expectations. Since the elements $\partial_{s}a_{k_{i_1}...i_2j_1...j_2}$ are forms with values in the Weyl algebra we look for a linear operator which transforms a $k$-form into a $(k + 1)$-form. This condition is fulfilled by the exterior covariant derivative operator [92].

**Definition 4.17.** The exterior covariant derivative $\partial : \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^k \rightarrow \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^{k+1}$ is a linear differential operator such that for every $a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda$

$$
\partial a := dq^s \wedge \partial_{s}a. \quad (4.1)
$$

Note that in Fedosov’s publications [71,72] the operator $\partial$ is simply called ‘connection’.

**Theorem 4.5.** [72] The exterior covariant derivative $\partial$ in the bundle $\mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda$ in a Darboux atlas $\{(U_\varphi, \varphi)\}_{\varphi \in \mathcal{I}}$ can be written as

$$
\partial a = da + \frac{1}{i\hbar} [\Gamma, a], \quad (4.2)
$$
where \( a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda \) and
\[
\Gamma := \frac{1}{2} \Gamma_{ij,k} dq^k
\] (4.3)
is a 1-form symmetric in \((i,j,k)\).

It is worthy to make the following remarks:

• The formula (4.2) is valid only in Darboux coordinates, because only in such charts \( \Gamma_{ijk} \) is symmetric in all the indices.

• We consider \( \Gamma \) like an element of \( \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^1 \) although we know that the elements \( \Gamma_{ijk} \) are not components of any tensor. This fact does not influence on the \( \circ \)-product in a Darboux chart.

• The formula (4.2) is similar to the definition of the connection matrix [90] but, in fact, the 1-form (4.3) is not the connection matrix. In the analyzed case the connection matrix is infinite dimensional.

**Theorem 4.6.** [72] The exterior covariant derivative preserves the degree of the forms of \( \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda \).

**Proof.** - The derivative \( d \) does not change the degree of the forms of \( \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda \). The commutator increases the power of \( \hbar \) in one unit since it has two derivatives in \( X^i \)'s, but the fact of dividing by \( i \hbar \) finally preserves the degree.

**Theorem 4.7.** For every two forms \( a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^{k_1} \) and \( b \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^{k_2} \) we have
\[
\partial(a \circ b) = \partial a \circ b + (-1)^{k_1} a \circ \partial b.
\] (4.4)

**Proof.** - We make the proof in Darboux coordinates.

\[
\partial(a \circ b) := d(a \circ b) + \frac{1}{i\hbar} [\Gamma, a \circ b]
\]
\[
= da \circ b + (-1)^{k_1} a \circ db + \frac{1}{i\hbar}(-1)^{1-k_1} a \circ [\Gamma, b] + \frac{1}{i\hbar} [\Gamma, a] \circ b = \partial a \circ b + (-1)^{k_1} a \circ \partial b.
\] (4.5)

Note that the relation (4.4) is a Leibniz rule for the exterior covariant derivative \( \partial \) and the \( \circ \)-product.

Let us consider now the second exterior covariant derivative \( \partial(\partial a) \). As before we make our computations in a Darboux chart. Thus,
\[
\partial(\partial a) = \partial \left( da + \frac{1}{i\hbar} [\Gamma, a] \right) = d(da) + d \left( \frac{1}{i\hbar} [\Gamma, a] \right) + \frac{1}{i\hbar} [\Gamma, da] + \frac{1}{i\hbar} [\Gamma, a].
\] (4.5)

Using
\[
d(a \wedge b) = da \wedge b + (-1)^{k_1} a \wedge db, \quad a \in \Lambda^k,
\]

\[19\]
which is a general property valid for any form, and remembering that \( \Gamma \) is a 1-form, we obtain
\[
d(\Gamma \circ a) = d\Gamma \circ a - \Gamma \circ da
\]
\[
d(a \circ \Gamma) = da \circ \Gamma + (-1)^k a \circ d\Gamma.
\](4.6)

From (3.11) and (4.6) we obtain
\[
d[\Gamma, a] = [d\Gamma, a] - [\Gamma, da].
\](4.7)

From the Jacobi identity (3.13)
\[
(1+1)(-1)^{(1+k)}[\Gamma, [\Gamma, a]] + (-1)^{(1+k)}[\Gamma, [\Gamma, \Gamma]] + (-1)^k(1+1)[\Gamma, [\Gamma, \Gamma]] = 0
\](4.8)

and using the relation
\[
[\Gamma, [a, \Gamma]] = (1+1)(-1)^{(1+k)}[\Gamma, [\Gamma, a]]
\](4.9)

we see that
\[
2[\Gamma, [\Gamma, a]] = -[a, [\Gamma, \Gamma]]
\](4.10)

Moreover,
\[
[\Gamma, \Gamma] = 2\Gamma \circ \Gamma.
\](4.11)

Putting (4.7) and (4.10) with (4.11) in (4.5) we get
\[
\partial(\partial a) = \frac{1}{i\hbar}[d\Gamma + \frac{1}{i\hbar}\Gamma \circ \Gamma, a].
\](4.12)

We see that the 2-form
\[
R = \Gamma + \frac{1}{i\hbar}\Gamma \circ \Gamma
\](4.13)

is the curvature of the connection \( \Gamma \). In Darboux charts we can write
\[
\partial(\partial a) = \frac{1}{i\hbar}[R, a].
\](4.14)

Let us find the explicit form of the 2-form of curvature (4.13). For simplicity we consider elements of the Weyl algebra bundle \( P^* \mathcal{M}[[\hbar]] \) acting on vectors. By \( X \) we denote an arbitrary fixed vector field on \( M \).

\[
d\Gamma(X, X) = d \left( \frac{1}{2} \Gamma_{ijk} X^i X^j dq^k \right) = \frac{1}{2} \frac{\partial \Gamma_{ijk}}{\partial q^l} X^i X^j dq^l \wedge dq^k
\]
\[
= \frac{1}{4} \left( \frac{\partial \Gamma_{ikl}}{\partial q^j} - \frac{\partial \Gamma_{jkl}}{\partial q^i} \right) X^i X^j dq^l \wedge dq^k = \frac{1}{4} \omega_{is} \left( \frac{\partial \Gamma_{ijl}}{\partial q^s} - \frac{\partial \Gamma_{jsl}}{\partial q^i} \right) X^i X^j dq^l \wedge dq^k,
\](4.15)

and
\[
\Gamma \circ \Gamma(X, X) := \frac{i\hbar}{2} \omega_{ij} \cdot 2 \cdot 2 \cdot \frac{1}{2} \Gamma_{ijkl} X^i dq^j \wedge dq^k
\]
\[
= \frac{i\hbar}{2} \omega_{ij} \cdot 2 \cdot 2 \cdot \frac{1}{2} \Gamma_{ijkl} X^i X^j dq^l \wedge dq^k = -\frac{i\hbar}{2} \Gamma_{ij}^{jk} \Gamma_{ijkl} X^i X^j dq^l \wedge dq^k
\]
\[
= -\frac{i\hbar}{2} \omega_{js} \cdot 2 \cdot 2 \cdot 2 \cdot \frac{1}{2} \Gamma_{ijkl} X^i X^j dq^l \wedge dq^k = \frac{i\hbar}{4} \omega_{is} \left( \Gamma_{jk}^{ul} \Gamma_{ul}^{rs} - \Gamma_{jl}^{ul} \Gamma_{ul}^{rs} \right) X^i X^j dq^l \wedge dq^k.
\](4.16)
From (4.15) and the last expression of (4.16) we obtain that
\[ R(X, X) = \frac{1}{4} R_{ijkl} X^i X^j dq^l \wedge dq^k \]  
where
\[ R_{ijkl} = \omega_{si} \left( \frac{\partial \Gamma^s_{kj}}{\partial q^l} - \frac{\partial \Gamma^s_{lj}}{\partial q^k} + \Gamma^u_{kj} \Gamma^s_{ul} - \Gamma^u_{lk} \Gamma^s_{ju} \right) \]
is the curvature tensor (2.22) of the symplectic connection \( \Gamma_{ijk} \) on the phase space \( M \).

From (4.17) we can see easily that \( R_{ijkl} \) is symmetric in the indices \( \{i, j\} \) and antisymmetric in \( \{k, l\} \).

It is can be proved that although we work in a Darboux chart, the relations (4.17) and (4.14) hold in every chart on \( M \).

Let us introduce two operators acting on the bundle \( \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda \). To make the notation more clear we will operate on elements of the Weyl algebra acting on vectors of the tangent space \( TM \) to the symplectic manifold \( M \), i.e. \( a_{i_1, \ldots, i_l, j_1, \ldots, j_s} dq^{i_1} \wedge \cdots \wedge dq^{i_l} X^{i_1} \cdots X^{i_l} \).

**Definition 4.18.** The operator \( \delta : \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^s \to \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^{s+1} \) defined as
\[ \delta a := dq^k \wedge \left( \frac{\partial a}{\partial X^k} \right) \]
is known as antiderivation.

In order to see if the above definition is consistent, let us compute the antiderivation of the constant tensor field. From the equation (4.19) we obtain
\[
\delta(1_{i_1, \ldots, i_l, j_1, \ldots, j_s} dq^{i_1} \wedge \cdots \wedge dq^{i_l}) = X^{i_1} \cdots X^{i_l} dq^{i_1} \wedge dq^{i_2} \wedge \cdots \wedge dq^{i_l} + X^{i_1} X^{i_2} \cdots X^{i_l} dq^{i_2} \wedge dq^{i_1} \wedge \cdots \wedge dq^{i_l} + \cdots + X^{i_1} \cdots X^{i_{l-1}} dq^{i_{l-1}} \wedge dq^{i_1} \wedge \cdots \wedge dq^{i_l}.
\]
Hence
\[
\delta(1_{i_1, \ldots, i_l, j_1, \ldots, j_s} dq^{i_1} \wedge \cdots \wedge dq^{i_l}) = 1_{i_2, \ldots, i_l, j_1, \ldots, j_s} dq^{i_1} \wedge dq^{i_2} \wedge \cdots \wedge dq^{i_l} + 1_{i_1, i_3, \ldots, i_l, j_1, \ldots, j_s} dq^{i_2} \wedge dq^{i_1} \wedge \cdots \wedge dq^{i_l} + \cdots + 1_{i_1, \ldots, i_{l-1}, j_1, \ldots, j_s} dq^{i_l} \wedge dq^{i_1} \wedge \cdots \wedge dq^{i_l},
\]
which, of course, belongs to \( \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^{s+1} \).

**Theorem 4.8.** The operator \( \delta \) may be written as a 1-form
\[ \delta = \frac{1}{i\hbar} [\omega_{ij} X^i dq^j, \cdot]. \]

**Theorem 4.9.** The operator \( \delta \) lowers the degree of the elements of \( \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda \) of 1.

The proofs of both theorems are straightforward.
Definition 4.19. The operator $\delta^* : \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^s \to \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^{s-1}$ is defined as

$$\delta^* a := X^k \left( \frac{\partial}{\partial q^k} \right) a.$$  \hspace{1cm} (4.21)

It can be considered as the ‘opposite’ of the antiderivation operator. The symbol $\langle \rangle$ denotes the inner product. Effectively, from the definition (4.21) we obtain that

$$\delta^* (X^{i_1} \ldots X^{i_n} dq^{j_1} \land \ldots \land dq^{j_s}) =$$

$$X^{j_1} X^{i_1} \ldots X^{i_n} dq^{j_2} \land \ldots \land dq^{j_s} - X^{j_2} X^{i_1} \ldots X^{i_n} dq^{j_1} \land dq^{j_3} \land \ldots \land dq^{j_s}$$

$$+ \ldots + (-1)^{s+1} X^{j_s} X^{i_1} \ldots X^{i_n} dq^{j_1} \land \ldots \land dq^{j_{s-1}}. \hspace{1cm} (4.22)$$

Theorem 4.10. [72] The operator $\delta^*$ raises the degree of the forms of $\mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda$ in 1.

Proof.- The operator $\delta^*$ exchanges the 1-form $dq^k$ into $X^k$ and does not touch $\hbar$, so the degree increases in 1.

Theorem 4.11. [72] The operators $\delta$ and $\delta^*$ do not depend on the choice of local coordinates and have the following properties:

(i) linearity,

(ii) $\delta^2 = (\delta^*)^2 = 0$,

(iii) for the monomial $X^{i_1} \ldots X^{i_n} dq^{j_1} \land \ldots \land dq^{j_s}$ we have

$$(\delta \delta^* + \delta^* \delta) X^{i_1} \ldots X^{i_n} dq^{j_1} \land \ldots \land dq^{j_s} = (l + s) X^{i_1} \ldots X^{i_n} dq^{j_1} \land \ldots \land dq^{j_s}.$$

Proof.- Linearity is obvious, thus let us start from the second property. From the symmetry of $X^{i_1} \ldots X^{i_n} dq^{j_1} \land \ldots \land dq^{j_s}$ in $X^i$’s it is sufficient to compute $\delta^2$ for two fixed $X^i$’s. Hence

$$\delta^2 (X^{i_1} X^{i_2}) = \delta (X^{i_2} dq^{i_1} + X^{i_1} dq^{i_2}) = dq^{i_2} \land dq^{i_1} + dq^{i_1} \land dq^{i_2} = 0.$$ 

Now from the antisymmetry of $X^{i_1} \ldots X^{i_n} dq^{j_1} \land \ldots \land dq^{j_s}$ in $dq^j$’s it is enough to find $(\delta^*)^2 (dq^{j_1} \land dq^{j_2})$ for two fixed $dq^{j_1}$ and $dq^{j_2}$. So,

$$(\delta^*)^2 (dq^{j_1} \land dq^{j_2}) = \delta^* (X^{j_1} dq^{j_2} - X^{j_2} dq^{j_1}) = X^{j_1} X^{j_2} - X^{j_2} X^{j_1} = 0.$$ 

To show the third statement we consider the monomial $X^i dq^j$. The result is easy to generalize on $X^{i_1} \ldots X^{i_n} dq^{j_1} \land \ldots \land dq^{j_s}$ using the symmetry in $X^i$’s and the antisymmetry in $dq^j$’s. Now

$$(\delta \delta^* + \delta^* \delta) (X^i dq^j) = \delta (X^i X^j) + \delta^* (dq^i \land dq^j) = X^i dq^j + X^j dq^i + X^i dq^j - X^j dq^i = 2X^i dq^j.$$ 

\hspace{1cm} \Box
Definition 4.20. [72] There is an operator $\delta^{-1} : \mathcal{P}^*\mathcal{M}[[\hbar]] \otimes \Lambda^s \to \mathcal{P}^*\mathcal{M}[[\hbar]] \otimes \Lambda^{s-1}$ defined by

$$\delta^{-1} a := \begin{cases} \frac{1}{l+s} \delta^* a & \text{for } l + s > 0 \\ 0 & \text{for } l + s = 0 \end{cases} \quad (4.23)$$

where $l$ is the degree of $a$ in $X^i$’s (i.e. the number of $X$’s) and $s$ is the degree of the form.

Note that, in fact, $\delta^{-1}$ is, up to constant, the $\delta^*$ operator.

The straightforward consequence of the linearity and the decomposition on monomials is the de Rham decomposition of the form $a \in \mathcal{P}^*\mathcal{M}[[\hbar]]\Lambda$ as shows next theorem.

Theorem 4.12. [72] For every $a \in \mathcal{P}^*\mathcal{M}[[\hbar]]\Lambda$ the following equality holds

$$a = \delta\delta^{-1} a + \delta^{-1}\delta a + a_{00}, \quad (4.24)$$

where $a_{00}$ is a function on the symplectic manifold $M$.

Proof.- For functions $a_{00}$ on $M$ the fact that the relation (4.24) holds is evident. To show that the de Rham decomposition is true also for elements of $\mathcal{P}^*\mathcal{M}[[\hbar]]\Lambda$ we use the fact that the operators $\delta$ and $\delta^{-1}$ are linear. From the theorem (4.11) we have for a monomial that

$$(\delta\delta^* + \delta^*\delta)(X^{i_1} \ldots X^{i_l} dq^{j_1} \wedge \ldots \wedge dq^{j_s}) = (l + s) X^{i_1} \ldots X^{i_l} dq^{j_1} \wedge \ldots \wedge dq^{j_s}. \quad (4.25)$$

From the definition (4.23) the l.h.s. of the above equation may be written as

$$(\delta(l + s)\delta^{-1} + ((l + 1) + (s - 1))\delta^{-1}\delta) X^{i_1} \ldots X^{i_l} dq^{j_1} \wedge \ldots \wedge dq^{j_s}. \quad (4.25)$$

Hence, from (4.25) we immediately obtain (4.24).

The definition (1.17) of the exterior covariant derivative is based on covariant derivatives determined by the symplectic connection. It is possible to generalize this description of exterior covariant derivative (compare with [91] where the term ‘connection’ is used instead of ‘exterior covariant derivative’).

Definition 4.21. The exterior covariant derivative $\tilde{\partial} : \mathcal{P}^*\mathcal{M}[[\hbar]] \otimes \Lambda^0 \to \mathcal{P}^*\mathcal{M}[[\hbar]] \otimes \Lambda^1$ is a linear differential operator such that for every $a \in \mathcal{P}^*\mathcal{M}[[\hbar]]\otimes \Lambda^0$ and $f \in \Lambda^0$

$$\tilde{\partial}(f \cdot a) = df \otimes a + f \cdot \tilde{\partial}a. \quad (4.26)$$

This definition may be extended to an arbitrary $k$-form with values in the Weyl algebra $\mathcal{P}^*\mathcal{M}[[\hbar]]$.

Theorem 4.13. [91] There is a unique operator $\tilde{\partial} : \mathcal{P}^*\mathcal{M}[[\hbar]] \otimes \Lambda^k \to \mathcal{P}^*\mathcal{M}[[\hbar]] \otimes \Lambda^{k+1}$ satisfying:

1. $$\tilde{\partial}(f \wedge a) = df \wedge a + (-1)^s f \wedge \tilde{\partial}a \quad (4.27)$$

for every $f \in \Lambda^s$, $a \in \mathcal{P}^*\mathcal{M}[[\hbar]] \otimes \Lambda^k$.  

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for $a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^0$.

In a local Darboux chart we can write

$$\check{\partial} a = da + \frac{1}{i\hbar} [\Phi, a],$$

(4.29)

where $\Phi \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^1$ is a connection determining the derivation $\check{\partial}$. The same definition works for 0-forms, so if $a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^0$

$$\check{\partial} a = da + \frac{1}{i\hbar} [\Phi, a].$$

(4.30)

Indeed, the operator $d + \frac{1}{i\hbar} [\Phi, \cdot]$ is linear. For every $a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda^0$ and $f \in \Lambda^0$

$$\check{\partial} a = df \cdot a + (-1)^s f \wedge da + f \cdot \frac{1}{i\hbar} [\Phi, a] = df \cdot a + f \cdot \check{\partial} a.$$  

Hence, the condition (4.26) is fulfilled. Now for $a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda^k$ and $f \in \Lambda^s$ we see from (4.29) that the property (4.27) holds, i.e.

$$\check{\partial} (f \wedge a) = d(f \wedge a) + \frac{1}{i\hbar} [\Phi, f \wedge a]$$

$$= df \wedge a + (-1)^s f \wedge da + \frac{1}{i\hbar} \Phi \circ (f \wedge a) - (-1)^{(s+k)} \frac{1}{i\hbar} (f \wedge a) \circ \Phi$$

$$= df \wedge a + (-1)^s f \wedge da + (-1)^s f \wedge \frac{1}{i\hbar} \Phi \circ a - (-1)^s f \wedge (-1)^k \frac{1}{i\hbar} a \circ \Phi$$

$$= df \wedge a + (-1)^s f \wedge (da + \frac{1}{i\hbar} [\Phi, a])$$

$$= df \wedge a + (-1)^s f \wedge \check{\partial} a.$$  

Since $\check{\partial} a = \check{\partial} a$ for every $a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \otimes \Lambda^0$ we can use the formula (4.29) as definition of the exterior covariant derivative in $\mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda$.

Analogously to the case of the connection $\Gamma$ (4.13) we define the curvature of $\Phi$ as a 2-form with values in the Weyl algebra

$$\check{\mathcal{R}} = d\Phi + \frac{1}{i\hbar} \Phi \circ \Phi.$$  

(4.31)

Of course, in Darboux coordinates for every $a \in \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda$ we can write

$$\check{\partial} (\check{\partial} a) = \frac{1}{i\hbar} [\check{\mathcal{R}}, a].$$  

(4.32)

Among the properties of the curvature $\check{\mathcal{R}}$ we present the so-called Bianchi identity.

**Theorem 4.14.** For the curvature $\check{\mathcal{R}}$ the following relation holds

$$\check{\partial} \check{\mathcal{R}} = 0.$$  

(4.33)
Proof.-

\[
\check{\partial} R = d\check{R} + \frac{1}{i\hbar} [\Phi, \check{R}] \quad \text{where} \quad d \left( d\Phi + \frac{1}{i\hbar} \Phi \circ \Phi \right) + \frac{1}{i\hbar} [\Phi, d\Phi + \frac{1}{i\hbar} \Phi \circ \Phi] \\
= d^2\Phi + \frac{1}{i\hbar} d\Phi \circ \Phi + (-1)^1 \frac{1}{i\hbar} d\Phi + \frac{1}{i\hbar} \Phi \circ d\Phi - \frac{1}{i\hbar} (-1)^2 d\Phi \circ \Phi + \frac{1}{i\hbar} \Phi = 0. 
\]

\[
\frac{1}{i\hbar} \Phi \circ \Phi - (-1)^1 \frac{1}{i\hbar} \Phi \circ \Phi = 0. 
\]

In our considerations a special role will be played by the so called ‘abelian connection’.

**Definition 4.22.** [72] A connection \( D \) (defined by (4.28)) in the bundle \( \mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda \) is called abelian if for any section \( a \in C^\infty(\mathcal{P}^* \mathcal{M}[[\hbar]] \Lambda) \)

\[
D^2 a = \frac{1}{i\hbar} [\Omega, a] = 0. 
\]

(4.34)

The curvature \( \Omega \) of the abelian connection \( D \) is a central form. Since \( \Omega \) is central we deduce that

\[
d\Omega = 0. 
\]

(4.35)

Indeed, from the general formula (4.29) we have

\[
D\Omega = d\Omega + \frac{1}{i\hbar} [\check{\Gamma}, \Omega], 
\]

where by \( \check{\Gamma} \) we denote the 1-form of the abelian connection. Since \( \Omega \) is a central form the commutator \([\check{\Gamma}, \Omega]\) disappears. From the other side, the Bianchi identity (4.33) holds for the curvature \( \Omega \), i.e. \( D\Omega = 0 \). Hence \( d\Omega = 0 \).

Let us assume that in a Darboux chart

\[
D = \partial + \frac{1}{i\hbar} [(\omega_{ij} X^i dq^j + r), \cdot], 
\]

(4.36)

where \( r \) is a globally defined 1-form satisfying the Weyl normalizing condition, i.e. the part of \( r \) not containing \( X^i \)'s vanishes. Let us compute the curvature of the connection (4.36). As we have mentioned before in (4.31)

\[
\Omega = d\check{\Gamma} + \frac{1}{i\hbar} \check{\Gamma} \circ \check{\Gamma}, 
\]

where

\[
\check{\Gamma} := \Gamma + \omega_{ij} X^i dq^j + r. 
\]

(4.37)

We can see that

\[
\Omega = d\Gamma + dr + \frac{1}{i\hbar} \Gamma \circ \Gamma + \frac{1}{i\hbar} \Gamma \circ \omega_{ij} X^i dq^j + \frac{1}{i\hbar} \Gamma \circ r + \frac{1}{i\hbar} \omega_{ij} X^i dq^j \circ \Gamma + \frac{1}{i\hbar} \omega_{ij} X^i dq^j \circ r + \frac{1}{i\hbar} \omega_{ij} X^i dq^j \circ \Gamma + \frac{1}{i\hbar} \omega_{ij} X^i dq^j \circ r. 
\]

(4.38)

We know that

\[
d\Gamma + \frac{1}{i\hbar} \Gamma \circ \Gamma = R. 
\]

(4.39)
Since our definition has been formulated in a Darboux chart we can write
\[ dr + \frac{1}{i\hbar} \Gamma \circ r + \frac{1}{i\hbar} r \circ \Gamma = dr + \frac{1}{i\hbar} [\Gamma, r] = \partial r. \] (4.40)

Moreover,
\[ \frac{1}{i\hbar} \Gamma \circ \omega_{ij} X^i dq^j + \frac{1}{i\hbar} \omega_{ij} X^i dq^j \circ \Gamma = \frac{1}{i\hbar} [\Gamma, \omega_{ij} X^i dq^j] = \frac{1}{2} \cdot 2 \cdot \frac{1}{2} \omega_l^l \Gamma_{tr} X^r dq^s \wedge \omega_{ij} dq^j. \] (4.41)

We know that (2.6) holds, hence (4.41) equals to
\[ \delta_j^l \Gamma_{tr} X^r dq^s \wedge dq^j = \Gamma_{tr} X^r dq^s \wedge dq^j = 0, \] (4.42)

because \( \Gamma_{tr} \) is symmetric in all the indices. Now
\[ \frac{1}{i\hbar} \omega_{ij} X^i dq^j \circ r + \frac{1}{i\hbar} r \circ \omega_{ij} X^i dq^j = \frac{1}{i\hbar} [\omega_{ij} X^i dq^j, r] \overset{4.20}{=} -\delta r. \] (4.43)

Finally
\[ \frac{1}{i\hbar} \omega_{ij} X^i dq^j \circ \omega_{kl} X^k dq^l = -\frac{1}{2} \omega_{ij} dq^i \wedge dq^j. \] (4.44)

Putting (4.39)–(4.44) into (4.38) we obtain
\[ \Omega = -\frac{1}{2} \omega_{ij} dq^i \wedge dq^j + R - \delta r + \partial r + \frac{1}{i\hbar} r \circ r. \] (4.45)

The abelian property will be fulfilled, provided
\[ \delta r = R + \partial r + \frac{1}{i\hbar} r \circ r. \] (4.46)

**Theorem 4.15.** [72] The equation (4.46) has a unique solution satisfying the following conditions
\[ \deg r \geq 3, \quad \delta^{-1} r = 0. \] (4.47)

**Proof.-** From the decomposition (4.24) we obtain
\[ r = \delta^{-1} \delta r, \]
because \( r \) is a 1-form we have \( r_{00} = 0 \) and \( \delta^{-1} r = 0 \). Moreover, \( \delta r \) is a solution of (4.46).
Hence
\[ r = \delta^{-1} (R + \partial r + \frac{1}{i\hbar} r \circ r). \] (4.48)

The operator \( \delta^{-1} \) raises the degree by 1, so (4.48) is a recurrent formula starting with \( \delta^{-1} R \). The proof that solution (4.48) really fulfills (4.46) and that it is unique, is more complicated (see it in [72] or [78]).

Note that we have not included in the expression of \( r \) terms with degrees lower than 3. This is due to the fact that the abelian connection \( \tilde{\Gamma} \) (4.37) contains terms with degrees 1 and 2 and for that \( r \), from its definition, is an object with degree \( \deg r \geq 3 \).
Theorem 4.16. For the abelian connection $D$ and two forms $a \in \mathcal{P} \Lambda^k \mathbb{M}[[\hbar]]$ and $b \in \mathcal{P} \Lambda^l \mathbb{M}[[\hbar]]$ we have
\[
D(a \circ b) = Da \circ b + (-1)^ka \circ Db.
\] (4.49)

The proof is analogous to that of theorem 4.7.

Corollary 4.3. [72] The set of 0-forms such that their abelian connection vanishes constitutes the subalgebra $\mathcal{P} \Lambda^0 \mathbb{M}[[\hbar]]$ of $\mathcal{P} \Lambda \mathbb{M}[[\hbar]]$.

5 The $\ast$-product on $M$

In former sections we have studied the structure of the Weyl algebra bundle. Thus, we are able to introduce the $\ast$-product on the symplectic manifold $M$.

Definition 5.23. The projection $\sigma : \mathcal{P} \Lambda \mathbb{M}[[\hbar]] \to C^\infty(M)$ assigns to each 0-form, $a$, of $\mathcal{P} \Lambda \mathbb{M}[[\hbar]]$ its part $a_{00}$ (according to its the de Rham decomposition (4.24)), i.e.
\[
\sigma(a) := a_{00}.
\] (5.1)

The following theorem holds.

Theorem 5.17. [72] For any function $a_{00} \in C^\infty(M)$ there exists a unique section $a$ of $C^\infty(\mathcal{P} \Lambda \mathbb{M}_D[[\hbar]])$ such that $\sigma(a) = a_{00}$. This element is defined by the recurrent formula
\[
a = a_{00} + \delta^{-1}(\partial a + \frac{1}{i\hbar}[r, a]).
\] (5.2)

The recurrent character of that solution can be proved analogously to (4.48) (for more details see [72]).

The map $\sigma$ gives a one-to-one correspondence between $C^\infty(\mathcal{P} \Lambda \mathbb{M}_D[[\hbar]])$ and $C^\infty(M)$. Then, $a = \sigma^{-1}(a_{00})$, where $\sigma^{-1}$ is the inverse map of $\sigma$. This map provides a quantization procedure as follows.

Definition 5.24. [72] Let $F_1, F_2$ be two $C^\infty(M)$-functions. The $\ast$-product is defined as
\[
F_1 \ast F_2 := \sigma(\sigma^{-1}(F_1) \circ \sigma^{-1}(F_2)).
\] (5.3)

This $\ast$-product can be considered as a generalization of the Moyal product of Weyl type defined for the case $M = \mathbb{R}^{2n}$. It has the following properties:

1. Invariance under Darboux transformations. In fact, the $\ast$-product is invariant under all smooth transformation of coordinates on $M$. However, the definitions of exterior covariant derivatives $\partial$ and $D$ cannot be written in the form (4.2) and (4.36), respectively.
2. In the limit $\hbar \to 0^+$ the $*$-product turns into the commutative point-wise multiplication of functions, i.e.

$$\lim_{\hbar \to 0^+} F_1 * F_2 = F_1 \cdot F_2. \quad (5.4)$$

Indeed, since in (5.3) only terms not containing $X^i$'s and $\hbar^n$, $n \geq 1$, are taking we can see that, in fact, only the terms of degree zero are essential in the product $\sigma^{-1}(F_1) \circ \sigma^{-1}(F_2)$. From the third property of the $\circ$-product (3.7) we deduce that the part of $\sigma^{-1}(F_1) \circ \sigma^{-1}(F_2)$ with degree zero is just $F_1 \cdot F_2$. Hence, we obtain (5.4).

3. The multiplication (5.3) is associative but noncommutative.

Effectively, from the theorem (5.17) the relation

$$\sigma^{-1}(F_1) = \sigma^{-1}(F_2) \iff F_1 = F_2$$

holds. It means that

$$\sigma^{-1}((F_1 * F_2) * F_3) = \sigma^{-1}(F_1 * F_2) \circ \sigma^{-1}(F_3) = (\sigma^{-1}(F_1) \circ \sigma^{-1}(F_2)) \circ \sigma^{-1}(F_3). \quad (5.5)$$

Since the $\circ$-product is associative the r.h.s. of (5.3) can be written as

$$\sigma^{-1}(F_1) \circ (\sigma^{-1}(F_2) \circ \sigma^{-1}(F_3)) = \sigma^{-1}((F_2 * F_3)) \circ \sigma^{-1}(F_1).$$

Hence

$$(F_1 * F_2) * F_3 = F_1 * (F_2 * F_3), \quad \forall F_1, F_2, F_3 \in C^\infty(M).$$

4. In the case $M = \mathbb{R}^{2n}$ the product defined above is just the Moyal product of Weyl type. Instead of proving this statement here we will analyze the case $M = \mathbb{R}^{2n}$ in the next section.

The properties mentioned above show that the $*$-product constructed according to the Fedosov idea is a natural generalization of the $*$-product in the trivial case when $M = \mathbb{R}^{2n}$ (Moyal-product). Several different properties of the $*$-product may be found in [74] or [89].

The construction, based on fibre bundle theory, of a $*$-product of Weyl type is finished. Remember that giving a symplectic space $M$ equipped with a symplectic connection a Weyl algebra bundle is constructed. Operating on flat sections of this bundle a non-commutative but associative product of observables on $M$ is defined. It seems to be the generalization of the Moyal product.

In this paper we have only considered $*$-products of the Weyl type. Other kinds of $*$-products, whose geometrical origin is the same that of the $*$-product of the Weyl type, were analyzed in [78].

The existence of many $*$-products on the same symplectic manifold is closely related to the equivalence problem of $*$-products. Two $*$-products $*_1$ and $*_2$ are said to be equivalent iff there exists a differential operator $\hat{T}$ such that for every two functions $F_1, F_2$ of $C^\infty(M)$, for which expressions appearing below have sense, the following relation holds

$$F_1 *_1 F_2 = \hat{T}^{-1}\left(\hat{T} F_1 *_2 \hat{T} F_2\right). \quad (5.6)$$
It has been proved in [93] that all the *-products on a symplectic manifold are equivalent to the Weyl type *-product constructed according to the Fedosov recipe. For more details about the equivalence problem of *-products see [93] and [94].

The definition of the *-product is not sufficient to study completely most of the physical problems. To solve the eigenvalue equation for an observable $O$ or to find its average value we must know how to define the states. There is not a general answer to this question. It seems that in the framework of the Fedosov formalism the states are described by functionals $W$ over some functions defined on the phase space of the system but, in fact, our knowledge about such objects is rather poor.

Topics connected with the representation problem of a quantum state on a phase space such as ‘traciality’ operation or closeness of the *-product can be studied, for instance, in [72,95,96].

6 Examples

In this section we present three systems where the Fedosov construction is accomplished. The results obtained in two of them are known from the ‘traditional’ quantum mechanics. The third one, recently developed by us [89], does not have its counterpart in the formalism of operators in a certain Hilbert space.

6.1 Cotangent bundle $T^*\mathbb{R}^n$

In this first example, we consider the simple case of a physical system whose phase space is the vector space $(\mathbb{R}^{2n}, \omega)$. This space is covered with one chart $(\mathbb{R}^{2n}, \varrho)$, in which the symplectic form takes its natural shape, i.e.

$$\omega = \sum_{i=1}^{n} dq^{n+i} \wedge dq^i.$$  (6.1)

We choose the symplectic connection in such a way that all the coefficients vanish, i.e. $\Gamma_{ijk} = 0$, $1 \leq i,j,k \leq 2n$. The symplectic curvature tensor $R_{ijkl}$ also vanishes.

The abelian connection (4.37) in the Weyl bundle is given by

$$\tilde{\Gamma} = \omega_{ij} X^i dq^j$$

and its curvature is a central form (4.45)

$$\Omega = -\frac{1}{2} \sum_{i=1}^{n} dq^{n+i} \wedge dq^i.$$  (4.45)

Hence, $r = 0$. It means that for every $F \in C^\infty(\mathbb{R}^{2n})$ according to (6.2) we can write

$$\sigma^{-1}(F) = F + \delta^{-1}(dF) = \sum_{m=0}^{\infty} \sum_{i_1, i_2, \ldots, i_m}^{m} \frac{1}{m!} \frac{\partial^m F}{\partial q^{i_1} \ldots \partial q^{i_m}} X^{i_1} \ldots X^{i_m}.$$  (6.2)
From (1.3), (3.6) and (6.1) we can see that

$$\sigma^{-1}(F_1) \circ \sigma^{-1}(F_2) = F_1(Q + X) \ast W F_2(Q + X), \quad (6.3)$$

where \( Q = (q^1, \ldots, q^{2n}) \), \( X = (X^1, \ldots, X^{2n}) \). The symbol \( \ast_W \) denotes the \( \ast \)-product of Weyl type on \( \mathbb{R}^{2n} \) defined by the formula (1.3). We can write the equality relation in (6.3) because \( \mathbb{R}^{2n} \) and \( T\mathbb{R}^{2n} \) are isomorphic.

Now from (5.3) we obtain that

$$F_1 \ast F_2 = \sigma(F_1(Q + X) \ast W F_2(Q + X)) = F_1(Q) \ast W F_2(Q). \quad (6.4)$$

We conclude that when the symplectic space is just \((\mathbb{R}^{2n}, \omega)\), the \( \ast \)-product computed using the Fedosov method is exactly the usual \( \ast \)-product of Weyl type.

### 6.2 Harmonic oscillator

Let us consider a system with phase space \((\mathbb{R}^{2n}, \omega)\), where \( \omega = dp \wedge dq \), and Hamiltonian

$$H = \frac{p^2}{2m} + \frac{mq^2}{2}.$$  

In a chart with coordinates \((p, q)\) all the coefficients of the connection \( \Gamma_{ijk}, i, j, k = 1, 2 \), vanish. In the new Darboux coordinates \((H, \phi)\), related with the old ones \((p, q)\) according to

$$q = \sqrt{\frac{2H}{m}} \sin \phi, \quad p = \sqrt{2mH} \cos \phi, \quad (6.5)$$

the symplectic form is rewritten as

$$\omega = dH \wedge d\phi.$$  

From the transformation rule for connections (2.11) we immediately obtain that the symplectic connection 1-form in coordinates \((H, \phi)\) is given by

$$\Gamma = \frac{1}{4H}X^2X^2d\phi + \frac{1}{2H}X^1X^2dH + HX^1X^1dH. \quad (6.6)$$

Obviously, its curvature vanishes. That implies that also \( r = 0 \) (see (4.48)) and the symplectic curvature in the Weyl algebra bundle is central (4.45)

$$\Omega = -\frac{1}{2}dH \wedge d\phi,$$

and the abelian connection is

$$\tilde{\Gamma} = X^1d\phi - X^2dH + \Gamma.$$  

Let us look for the eigenvalues and eigenfunctions of the Hamiltonian \( H \) in coordinates \((H, \phi)\). The eigenvalue equation takes the form

$$H \ast W_E(H, \phi) = E \cdot W_E(H, \phi), \quad (6.7)$$
where $W_E(H, \phi)$ is a functional representation of an eigenstate with eigenvalue $E$. From (5.2) we can write
\[
\sigma^{-1}(H) = H + X^2 + HX^1X^1 + \frac{1}{4H}X^2X^2
\]
and
\[
\sigma^{-1}(W_E) = W_E + \frac{\partial W_E}{\partial H}X^1 + \frac{\partial W_E}{\partial \phi}X^2 + H\frac{\partial W_E}{\partial H}X^1X^1 + \frac{1}{2}\frac{\partial^2 W_E}{\partial \phi^2}X^1X^2 + \frac{1}{2}\frac{\partial^2 W_E}{\partial \phi^2}X^1X^1 + \ldots
\]
The series $\sigma^{-1}(W_E)$ is infinite but only the terms of $\text{deg} \leq 2$ are essential because $\sigma^{-1}(H)$ has degree 2. Computing $\sigma^{-1}(H) \circ \sigma^{-1}(W_E)$ and projecting the product on the phase space we finally obtain that
\[
H \ast W_E(H, \phi) = (H - E)W_E - \frac{\hbar^2}{4} \frac{\partial W_E}{\partial H} H - \frac{\hbar^2}{4} H \frac{\partial^2 W_E}{\partial H^2} - \frac{\hbar}{2} \frac{\partial W_E}{\partial \phi} - \frac{\hbar^2}{16H} \frac{\partial^2 W_E}{\partial \phi^2} = 0. \quad (6.9)
\]
Since the function $W_E$ is real we see that
\[
\frac{i\hbar}{2} \frac{\partial W_E}{\partial \phi} = 0. \quad (6.10)
\]
Thus, $W_E$ depends only on $H$ and may contain $\hbar$ in the denominator. It is possible that for some solutions of (6.9) the series $\sigma^{-1}(W_E)$ is not well defined. Such functions are not admissible because for them the product $\sigma^{-1}(H) \circ \sigma^{-1}(W_E)$ does not exist. However, it happens that admissible solutions of (6.9) are Wigner functions
\[
W_{E_n}(H) = \frac{1}{\pi \hbar} (-1)^n \exp \left( \frac{-2H}{\hbar} \right) L_n \left( \frac{4H}{\hbar} \right), \quad (6.11)
\]
where $L_n$ is a Laguerre polynomial, and the energy $E$ is quantized
\[
E_n = \hbar (n + \frac{1}{2}), \quad n = 0, 1, 2, \ldots \quad (6.12)
\]
This result is well known from traditional quantum mechanics. It can be also obtained in terms of the Moyal product (1.4) [12].

### 6.3 2-Dimensional phase space with constant curvature tensor

Let us consider the phase space $(\mathbb{R}^2, \omega)$. Topologically this phase space is homeomorphic to $\mathbb{R}^2$. We cover it with an atlas containing only one chart $(\mathbb{R}^2, \phi)$. In this chart the symplectic form is, as usual, $\omega = dp \wedge dq$, where $q$ denotes the spatial coordinate and $p$ the momentum conjugate to $q$. We assume that in the chart $(\mathbb{R}^2, \phi)$ the connection 1-form is
\[
\Gamma = \frac{1}{2} pX^1X^1 dq. \quad (6.13)
\]
The connection (6.13) is well defined globally since we cover the symplectic manifold \((\mathbb{R}^2, \omega)\) with only one chart. The curvature 2-form is

\[
R = -\frac{1}{2} X^1 X^1 dq \wedge dp. \quad (6.14)
\]

Hence, in the chart \((\mathbb{R}^2, \varrho)\) the curvature tensor has only one nonvanishing component which is constant

\[
R_{112} = -R_{121} = -\frac{1}{4}. \quad (6.15)
\]

The Ricci tensor is

\[
K_{11} = \frac{1}{4}. \quad (6.16)
\]

Since the connection 1-form is defined by the expression (6.13), we are able to build the abelian connection. Let us compute the series \(r\) defined by (4.48). The first term in the recurrent expression (4.48) is

\[
\delta^{-1} R = \frac{1}{8} X^1 X^1 X^2 dq - \frac{1}{8} X^1 X^1 X^1 dp. \quad (6.17)
\]

After some computations (for details see [89]) we obtain that the abelian connection is

\[
\tilde{\Gamma} = -X^1 dp + X^2 dq + \frac{1}{2} p X^1 X^1 dp - \frac{1}{8} X^1 X^1 X^1 dp + \frac{1}{8} X^1 X^1 X^2 dq
\]

\[
+ \frac{1}{128} X^1 X^1 X^1 X^1 dp - \frac{1}{128} X^1 X^1 X^1 X^1 X^2 dq
\]

\[
- \frac{1}{1024} X^1 X^1 X^1 X^1 X^1 X^1 X^1 X^1 dp + \frac{1}{1024} X^1 X^1 X^1 X^1 X^1 X^1 X^1 X^2 dq - \cdots \quad (6.18)
\]

Note that the abelian connection is an infinite series, the Planck constant \(\hbar\) does not appear in any term of the series, and in the series \(r\) (4.37) only two kind of terms are present:

\[
(-1)^m a_{2m+1} (X^1)^{2m} X^2 dq
\]

and

\[
(-1)^m a_{2m+1} (X^1)^{2m+1} dp, \quad m \in \mathbb{N} - \{0\},
\]

where \(a_{2m+1}\) is related to the so called Catalan numbers and is given by

\[
a_{2m+1} = \frac{2}{m} \left( \frac{2m-2}{m-1} \right) \frac{1}{16^m}, \quad m \geq 1. \quad (6.19)
\]

The expression (6.18) allows us to write eigenvalue equations for observables. Let us start with constructing the explicit form of the eigenvalue equation for momentum \(p\). The general form of this equation is

\[
p \ast W_p(q, p) = p \cdot W_p(q, p), \quad (6.20)
\]

where \(p\) denotes the eigenvalue of \(p\) and \(W_p(q, p)\) is the Wigner function associated to the eigenvalue \(p\). We find that

\[
\sigma^{-1}(p) = p + X^2 + \frac{1}{2} p X^1 X^1 + \frac{1}{8} X^1 X^1 X^2
\]

\[
- \frac{1}{128} X^1 X^1 X^1 X^1 X^2 + \frac{1}{1024} X^1 X^1 X^1 X^1 X^1 X^1 X^1 X^2 - \cdots \quad (6.21)
\]

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Every coefficient $b_{2m+1}$ appearing in the term $b_{2m+1}X_1^1 \cdots X_1^1X_2^2$, for $m > 0$, can be expressed in the form

$$b_{2m+1} = (-1)^{m+1}a_{2m+1},$$  \hspace{1cm} (6.22)

where $a_{2m+1}$ is defined by (6.19).

Much more complicated is to find the general formula of the series $\sigma^{-1}(W_p)$ representing the Wigner function $W_p(q,p)$. After long considerations we conclude that the function $W_p(q,p)$ depends only on $p$ and that the eigenvalue equation for $W_p(p)$ is the infinite differential equation

$$p \cdot \left( W_p(p) - \frac{1}{8} \hbar^2 \frac{d^2 W_p(p)}{dp^2} - \frac{1}{128} \hbar^4 \frac{d^4 W_p(p)}{dp^4} - \cdots \right) = p \cdot W_p(p).$$ \hspace{1cm} (6.23)

The eigenvalue equation for momentum $p$ is a differential equation of infinite degree. There is no general method for solving such equations. Therefore, we decided to look for the solution of the eigenvalue equation of $p^2$. As $p \cdot p = p \star p$ the Wigner function fulfilling the equation (6.20) satisfies also the relation

$$p^2 \cdot W_p(p) = p^2 W_p(p).$$ \hspace{1cm} (6.24)

We can see that the above relation is the modified Bessel equation

$$\frac{1}{4} \hbar^2 p^2 \frac{d^2 W_p(p)}{dp^2} + \frac{1}{4} \hbar^2 p \frac{dW_p(p)}{dp} + (p^2 - p^2) W_p(p) = 0.$$ \hspace{1cm} (6.25)

The general solution of (6.25) is a linear combination of the following form

$$W_p(p) = A \cdot I_{2p} \left( \frac{2p}{\hbar} \right) + B \cdot K_{2p} \left( \frac{2p}{\hbar} \right),$$ \hspace{1cm} (6.26)

with $I_{2p} \left( \frac{2p}{\hbar} \right)$ the modified Bessel function with complex parameter $\frac{2p}{\hbar}$ and $K_{2p} \left( \frac{2p}{\hbar} \right)$ the modified Bessel function of second kind with parameter $\frac{2p}{\hbar}$. This solution is defined for arguments $\frac{2p}{\hbar} > 0$ (see [97]). Note that the deformation parameter $\hbar$ appears in the denominator of the argument. We must be very careful because this fact may cause the nonexistence of the series $\sigma^{-1}(W_p(p))$.

The function $A \cdot I_{2p} \left( \frac{2p}{\hbar} \right)$ is complex. Its real part grows up to infinity for $x \to \infty$ and, hence, it is not normalizable. The imaginary part of $I_{2p} \left( \frac{2p}{\hbar} \right)$ is proportional to $K_{2p} \left( \frac{2p}{\hbar} \right)$. Thus, the only physically admissible solution of (6.25) is

$$W_p(p) = B \cdot K_{2p} \left( \frac{2p}{\hbar} \right).$$ \hspace{1cm} (6.27)

However, solutions defined on the whole $\mathbb{R}$ are required. It is impossible to define the solution of the equation (6.25) on the whole axis. The problem is that the modified Bessel function of second kind is not defined for the argument value $p = 0$. Moreover, $\lim_{p \to 0^+} K_{2p} \left( \frac{2p}{\hbar} \right)$ does not exist.
Let us assume that the Wigner function $W_p(p)$ is a generalized function over the Schwartz space $S(p)$ of test smooth functions tending to 0 when $p \to \pm\infty$ faster than the inverse of any polynomial.

We define the Wigner function $W_p(p)$ as:

1. $p < 0$

$$W_p(p) = \begin{cases} \frac{4}{\pi\hbar} \cosh \frac{p}{\hbar} K^{\frac{2p}{\hbar}}(-2p) & \text{for } p < 0 \\ 0 & \text{for } p \geq 0 \end{cases}$$ (6.28)

2. $p > 0$

$$W_p(p) = \begin{cases} 0 & \text{for } p \leq 0 \\ \frac{4}{\pi\hbar} \cosh \frac{p}{\hbar} K^{\frac{2p}{\hbar}}(2p) & \text{for } p > 0 \end{cases}$$ (6.29)

3. $p = 0$

$$W_0(p) = \begin{cases} \frac{2}{\pi\hbar} K_0(-2p) & \text{for } p < 0 \\ \frac{2}{\pi\hbar} K_0(2p) & \text{for } p > 0 \end{cases}$$ (6.30)

The example analyzed before is a particular choice of $G$ in the 2-form of curvature

$$R = \pm G^2 X^1 X^1 dq \wedge dp,$$ (6.31)

where $G$ is some positive constant. The solutions of the equation (6.31) are divided in two classes for $R = G^2 X^1 X^1 dq \wedge dp$ and $R = -G^2 X^1 X^1 dq \wedge dp$, respectively. For $G \to 0$ both of them tend to $\delta(p - \mathbf{p})$.

The eigenvalue equation for position $q$ takes the form

$$q \ast W_q(q,p) = q W_q(q,p),$$ (6.32)

where $q$ denotes the eigenvalue of the position $q$. It can be separated in two parts: the real part

$$q \cdot W_q(q,p) = q W_q(q,p)$$ (6.33)

and the imaginary part

$$\frac{1}{2\hbar} \frac{\partial W_q(q,p)}{\partial p} + \frac{1}{96\hbar^3} \frac{\partial^3 W_q(q,p)}{\partial p^3} + \cdots = 0.$$ (6.34)

Let us start with the equation (6.31). Multiplying it by $\frac{2}{\hbar}$, introducing a new variable $z = \frac{\sqrt{2p}}{\hbar}$ and defining $w_q(q, z) = \frac{\partial W_q(q, z)}{\partial z}$, we obtain the formula

$$w_q(q, p) + \frac{1}{24} \frac{\partial^2 w_q(q, p)}{\partial z^2} + \cdots = 0.$$ (6.35)

This is, in fact, a homogeneous linear differential equation of infinite degree. Its solution neither depends on the parameter $q$ nor the factor $G$. Since the $\lim_{G \to 0^+} w_q(q, z)$ must be 0 the only admissible solution of (6.35) is $w_q(q, z) = 0$ for every $q$. We see that the
Wigner eigenfunction $W_q(q, z)$ depends only on $q$. From (6.33) we immediately obtain that

$$W_q(q, z) = \delta(q - q), \quad \forall q.$$  \hspace{1cm} (6.36)

Similar considerations can be done for any arbitrary curvature 2-form $R$ of the form (6.31). The final result will be the same, i.e. the eigenfunctions $W_q(q)$ depend only on $q$ and have the same form like in the case of the flat space $\mathbb{R}^{2n}$ with $\Gamma = 0$.

It is well known [93,94] that on a 2-dimensional symplectic manifold all the $\ast$-products are equivalent. Thus, for every two products $\ast_1$ and $\ast_2$ the relation (5.6) holds. It means that the $\ast$-product considered in this example is equivalent to the Moyal product (1.4). That is true but we have not algorithm which enables us to construct the operator $\hat{T}$ from the equation (5.6). We cannot transform Wigner eigenfunctions and eigenvalues solutions of the eigenvalue equation

$$F(p, q) \ast_1 W_{F1}(p, q) = F_1 \cdot W_{F1}(p, q)$$

into eigenvalues and eigenfunctions of

$$F(p, q) \ast_2 W_{F2}(p, q) = F_2 \cdot W_{F2}(p, q)$$

though we are aware of the existence of some relation between them.

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**References**

[1] P. A. M. Dirac, *The Principles of Quantum Mechanics*, Oxford Univ. Press, Oxford 1958.

[2] K. B. Wolf, “The Heisenberg-Weyl ring in Quantum Mechanics” in *Group Theory and Its Applications*, vol III. Ed. E. Loebl, Academic Press, New York 1975.

[3] L. Cohen, *J. Math. Phys.* 7, 781 (1966).

[4] G.S. Agarwal, E. Wolf, *Phys. Rev. D* 2, 2161 (1970); *Phys. Rev. D* 2, 2206 (1970).

[5] J. Underhill, *J. Math. Phys.* 19, 1933 (1978).
[6] L. van Hove, Mem. de l’Acad. Royal de Belgique (Class des Sci.) XXVI, 61 (1951).
[7] N. Woodhouse, Geometric Quantization, Oxford Univ. Press, New York 1980.
[8] J. E. Moyal, Proc. Camb. Phil. Soc. 45, 99 (1949).
[9] H. Weyl, The theory of groups and quantum mechanics, Dover, New York 1931.
[10] E. P. Wigner, Phys. Rev. 40, 749 (1932).
[11] H. J. Groenewold, Physica 12, 405 (1946).
[12] F. Bayen, M.Flato, C. Fronsdal, A. Lichnerowicz and D. Sternheimer, Ann. Phys. NY 111, 61 (1978); Ann. Phys. NY 111, 111 (1978).
[13] J. M. Gracia-Bondía and J. C. Varilly, J. Math. Phys. 29, 869 (1988); J. Math. Phys. 29, 880 (1988).
[14] J. Tosiek and M. Przanowski, Acta Phys. Pol. B26, 1703 (1995).
[15] M. Gerstenhaber, Ann. of Math. 79, 59 (1964); Ann. of Math. 84, 1 (1966).
[16] V. Chari and A. Pressley, A guide to Quantum Groups, Cambridge Univ. Press, Cambridge 1994.
[17] J. P. Amiet, P. Hugenin, Mécaniques classique et quantique dans l’espace de phase, Université de Neuchatel, 1981.
[18] F. Antonsen, gr-qc/9712012, quant-ph/9710021, quant-ph/9802031
[19] D. Arnal, J. C. Cortet, P. Molin and C. Pinczon, J. Math. Phys. 24, 276 (1983).
[20] A. Ballesteros, M. Gadella and M. A. del Olmo, J. Math. Phys. 33, 3379 (1992).
[21] M. Bordemann, H. Herbig and S. Waldmann, mathQA/9901015.
[22] M. Bordemann, N. Neumaier, M. J. Pflaum and S. Waldmann, math/9811055
[23] M. Bordemann and S. Waldmann, Comm. Math. Phys. 195, 549 (1998).
[24] J. F. Cariñena, J. M. Gracia-Bondía and J. C. Varilly, J. Phys. A: Math. Gen. 23, 901 (1990).
[25] J. F. Cariñena and J. Clemente-Gallardo, Czech. J. of Phys. 48, 1307 (1998).
[26] P. Carruthers and F. J. Zachariasen, Rev. Mod. Phys. 55, 245 (1983).
[27] A. S. Cattaneo and G. Felder, math QA/0012228.
[28] J. Dito, Lett. Math. Phys. 20, 125 (1990).
[29] J. Dito, J. Math. Phys. 33, 791 (1992).
[30] R. Estrada, J. M. Gracia-Bondía and J. C. Varilly, *J. Math. Phys.* **30**, 2789 (1989).
[31] H. Figueroa, J. M. Gracia-Bondía and J. C. Varilly, *J. Math. Phys.* **31**, 2664 (1990).
[32] R. Fiorese and M. A. Lledó, math/QA 0111132.
[33] C. Fronsdal, *Rep. Math. Phys.* **15**, 111 (1978).
[34] M. Gadella, M.A. Martín, L. M. Nieto and M. A. del Olmo, *J. Math. Phys.* **32**, 1182 (1991).
[35] M.A. Martín and M.A. del Olmo, *J. Phys. A: Math. Gen.* **29**, 689 (1996).
[36] O. Arratia, M.A. Martín and M.A. del Olmo, *Rend. Sem. Mat. Univ. Pol. Torino* **54**, 191 (1996).
[37] O. Arratia and M.A. del Olmo, *Int. J. of Modern Phys. A*, **12**, 125 (1997); *Fortschritte der Physik*, **45**, 103 (1997).
[38] O. Arratia and M.A. del Olmo, *Reports on Math. Phys.* **40**, 149 (1997).
[39] O. Arratia, M.A. Martín and M.A. del Olmo, *Rev. R. Acad. Cienc. Serie A, Math.* **96**, 63 (2002).
[40] J. F. Plebański, M. Przanowski, J. Tosiek and F. J. Turrubiates, *Acta Phys. Pol. B* **31**, 561 (2000).
[41] J. A. González, M. A. del Olmo and J. Tosiek, *J. Optics B: Quantum Semiclass. Opt.* **5**, S306 (2003).
[42] M. Gadella, L. M. Nieto, J. M. Noriega and E. Santos, *J. Math. Phys.* **28**, 2961 (1987).
[43] M. Gadella and E. Santos, *J. Math. Phys.* **22**, 1651 (1981).
[44] H. García-Compeán, J. F. Plebański, M. Przanowski and F. J. Turrubiates, *Int. J. of Modern Phys. A* **16**, 2533 (2001).
[45] J. M. Gracia-Bondía, *Phys. Rev. A* **30**, 691 (1983).
[46] J. M. Gracia-Bondía and J. C. Varilly, *J. Phys. A: Math. Gen.* **21**, 879 (1988).
[47] J. M. Gracia-Bondía and J. C. Varilly, *J. Math. Phys.* **36**, 2691 (1995).
[48] J. M. Gracia-Bondía, *Generalized Moyal Quantization on Homogeneous Symplectic Spaces*, Contemporary Mathematics (1991).
[49] S. Gutt, *Lett. Math. Phys.* **7**, 249 (1983).
[50] T. Hakılığlu and A. J. Dragt, quant-ph/0108081
[51] E. Hawkins, math/9811049
[52] M. Hillery, R.F. O’Connell, M.O. Scully and E.P. Wigner, *Phys. Rep.* **106**, 121 (1984).

[53] M. Kontsevich, [q-alg/9709040](http://arxiv.org/abs/q-alg/9709040).

[54] M. Kontsevich, AG/0106006.

[55] J. Krüger and A. Poffyn, *Physica* **85A**, 84 (1976).

[56] N. Landsman, *J. Geom. and Phys.* **12**, 93 (1993).

[57] L. M. Nieto, *J. Phys. A: Math. Gen.* **24**, 1667 (1980).

[58] H. Omori, Y. Maeda and A. Yoshioka, *Adv. in Math.* **85**, 224 (1991).

[59] J. F. Plebański, M. Przanowski and J. Tosiek, *Acta Phys. Pol.*, **B27**, 1961 (1996).

[60] P. de M. Rios and G. M. Tuynman, [math-ph/0201044](http://arxiv.org/abs/math-ph/0201044).

[61] W.I. Tatarskij, *Usp. Fiz. Nauk* **139**, 587 (1983).

[62] D. Tamarkin, [q-alg/9705007](http://arxiv.org/abs/q-alg/9705007).

[63] J. C. Varilly and J. M. Gracia-Bondía, *J. Math. Phys.* **29**, 869 (1988).

[64] J. Varilly, *The Stratonovich-Weyl Correspondence: a general approach to Wigner functions*, BIBOS 345 (1988).

[65] J. C. Varilly and J. M. Gracia-Bondía, *Ann. Phys.* **190**, 107 (1989).

[66] C. Zachos, [hep-th/0110114](http://arxiv.org/abs/hep-th/0110114).

[67] A. Lichnerowicz, “Differential Geometry and Deformations” in *Symplectic Geometry*. Ed. A. Crumeyrolle and J. Grifone, Pitman, Boston 1983.

[68] M. Gadella, *Fortsch. Phys.* **43**, 229 (1995).

[69] D. Sternheimer, QA/9809056.

[70] G. Dito and D. Sternheimer, QA/0201168.

[71] B. Fedosov, *J. Diff. Geom.* **40**, 213 (1994).

[72] B. Fedosov, *Deformation Quantization and Index Theory*, Akademie, Berlin 1996.

[73] N. Steenrod, *The topology of fibre bundles*, Princeton Univ. Press, Princeton 1951.

[74] M. Bordemann and S. Waldmann, *Lett. Math. Phys.* **41**, 243 (1997).

[75] M. Bordemann, N. Neumaier and S. Waldmann, [q-alg/9711016](http://arxiv.org/abs/q-alg/9711016).

[76] M. Bordemann, [q-alg/9605038](http://arxiv.org/abs/q-alg/9605038).
[77] A. S. Cattaneo, G. Felder and L. Tomassini, math QA/0111290.
[78] M. Przanowski and J. Tosiek, Acta Physica Pol. B30, 179 (1999).
[79] A. Cannas de Silva, Lectures on Symplectic Geometry, Springer, Berlin 2001.
[80] P. Libermann, C. M. Marle, Symplectic Geometry and Analytical Mechanics, D. Reidel, New York 1986.
[81] V. Guillemin, S. Sternberg, Symplectic Techniques in Physics, Cambridge Univ. Press, Cambridge 1982.
[82] R. Abraham, J. E. Marsden, Foundations of Mechanics, Benjamin Cummings, Reading, 1978.
[83] V. I. Arnold, Mathematical Methods of Classical Mechanics, Springer, New York 1978.
[84] A. A. Kirillov, Elements of the Theory of Representations, Springer, Berlin 1976.
[85] W. D. Curtis and F. R. Miller, Differential Manifolds and Theoretical Physics, Academic Press, Orlando 1985.
[86] I. Vaisman, Monathefte fuer Mathematik 100, 299 (1985).
[87] I. Gel’fand, V. Retakh, M. Shubin, Fedosov Manifolds, presented at Symplectic Geometry Workshop, Toronto, June 1997.
[88] M. Nakahara, Geometry, Topology and Physics, Adam Hilger, Bristol 1990.
[89] M. Gadella, M. A. del Olmo and J. Tosiek, Ann. Phys. 307, 272 (2003).
[90] S. S. Chern, W. H. Chen, K. S. Lam, Lectures on Differential Geometry, World Scientific, Singapore 1999.
[91] R. S. Ward and R. O. Wells Jr., Twistor Geometry and Field Theory, Cambridge Univ. Press, Cambridge 1989.
[92] W. A. Poor, Differential Geometric Structures, McGraw-Hill, New York.
[93] M. Bertelson, M. Cahen and S. Gutt, Class. Quantum Grav. 14, A93 (1997).
[94] S. Gutt and J. Rawnsley, in Proceedings of the Workshop on Quantization and Momentum Maps at the University of Warwick 1997.
[95] H. Omori, Y. Maeda and A. Yoshioka, Lett. Math. Phys. 26, 285 (1992).
[96] A. Connes, M. Flato and D. Sternheimer, Lett. Math. Phys. 2, 1 (1992).
[97] G. N. Watson, A Treatise on the Theory of Bessel Functions, Cambridge Univ. Press, Cambridge 1958.