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Diagonal Representation for a Generic Matrix Valued Quantum Hamiltonian

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A general method to derive the diagonal representation for a generic matrix valued quantum Hamiltonian is proposed. This method can be applied to any kind of quantum system which has an energy band spectrum. In particular it would be useful for the study of relativistic Dirac particles in strong external fields as well as for Bloch electrons in solids. In this approach the diagonal energy operator and the dynamical operators which depend on Berry phase terms can be expanded in power series in the Planck constant. The connection with previous less general approaches is also presented. As an example, the case of a Dirac electron in an external electric field is considered.

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I. INTRODUCTION

The physical content of quantum systems is most often best revealed in the representation where the Hamiltonian is diagonal. The paradigmatic example is provided by the Foldy-Wouthuysen (FW) representation of the Dirac Hamiltonian for relativistic particles interacting with an external electromagnetic field. In this representation the positive and negative energy states are separately represented and the non-relativistic Pauli-Hamiltonian is obtained [1]. Actually even if several exact FW transformations have been found for some definite classes of potentials [2][3][4], the diagonalization (block-diagonalization for Dirac Hamiltonian) is a difficult mathematical problem requiring some approximations, essentially a perturbation expansion in weak fields. To overcome this limitation we have recently proposed a new method based on a expansion in powers of the Planck constant \hbar which

is not restricted to Dirac Hamiltonians (and thus very different from the FW method) but is applicable to a large class of quantum systems: as a non-Dirac example we carried out the Hamiltonian diagonalization of a spinless electron in magnetic Bloch bands [5][6] [7]. Another advantage of the \hbar expansion, apart from the fact that it is valid for strong fields, is that the semiclassical limit of (relativistic) quantum mechanics is readily obtained. Since semiclassical methods play a very important role in solid state physics too, in studying the dynamics of electrons to account for the various properties of metals, semiconductors and insulators, it is clear that a diagonalization method valid for a generic Hamiltonian is of a great interest. Actually, one must underline that there exists several other approaches treating the approximate diagonalization of general quantum Hamiltonians. The general method developed by Blount is based on an expansion in weak electromagnetic fields and is in principle applicable to Bloch electrons and Dirac electrons as well [8]. But his method is limited by construction to weak external perturbations. Moreover, it leads to a formal series expansion which leads unfortunately to cumbersome calculations for practical applications. Later on, Weigert and Littlejohn developed a systematic method to diagonalize general quantum Hamiltonian in a series expansion in \hbar [9]. Although similar in spirit, their method is very different from ours and more closely connected to Blount's one. It leads also to a formal series expansion written in terms of symbols of operators which also makes the method very complicated for practical applications. It is worth mentioning that recently a variant of the Foldy Wouthuysen transformation valid for strong fields and based also on an expansion in \hbar of the Dirac Hamiltonian was presented [10].

The method developed in [5] is very different from the ones previously discussed as it is based on a differential equation of the required diagonal Hamiltonian with respect to \hbar which is thus considered formally as a running parameter. It turns out that the resolution of this differential equation can be done by a series expansion in \hbar , giving thus, at least in principle, the diagonal Hamiltonian as a series in \hbar . Particularly interesting is the fact that the diagonal Hamiltonian in the semiclassical limit, which is often enough to get physical insight to the problem considered, is obtained by a straightforward integration of this differential equation. We could give in this case a general formula for an arbitrary matrix valued Hamiltonian in terms of covariant operators and commutators between Berry connections. This general formula was first derived by a direct diagonalization in [6] and allows us to study the dynamics of electrons and photons in a static gravitational field for which we

could go beyond other previous approaches [11]. Therefore, an other main advantage of the diagonalization procedure of [5] is that it embraces several different physical systems ranging from Bloch electrons in solid to Dirac particles interacting with any type of external fields even the strong ones. In [5], we also provided two examples in particle and solid state physics for which we could extend the diagonalization to the second order in \hbar , but due to the complexity of the problem we could not find a general expression to the second order in \hbar for a generic Hamiltonian.

In the present paper we solve this problem by presenting an alternative powerful method for the diagonalization of a generic matrix valued Hamiltonian that leads to a particularly compact and elegant exact expression for the required diagonal energy operator. The philosophy behind the approach developed in this paper, consists in a mapping of the initial quantum system to a classical one which can be diagonalized and then to return to the full quantum system. This method requires the introduction of some new mathematical objects like non-commuting operators evolving with the Planck constant promoted as a running variable. This construction leads us to define a differential calculus on a non-commutative space which shows some similarities with the stochastic calculus. It is not our goal in this paper to deepen the mathematical aspect of our approach that might look sometimes formal but also shows to be very elegant and much more powerful than any other ones as it is general and more tractable for practical applications.

The method considers \hbar as a running parameter as in [5] but is actually much simpler as it is not based on a differential equation. It leads directly to a formal exact diagonal representation for any arbitrary matrix valued Hamiltonian, which as we will prove in this paper, is an exact solution of the differential equation. It appears also, as one could expect, that the expansion in a series of \hbar of the exact diagonal Hamiltonian derived here is much more easier to obtain than by the successive integration of the differential equation. Consequently, contrary to the procedure in [5], one could find a general diagonal expression for any arbitrary Hamiltonian to the second order in \hbar (to higher order expressions becomes again very cumbersome but could in principle be deduced). As a physical illustration of this result we derive the full relativistic energy of a spin 1/2 particle in an external electric field to the second order in \hbar which allows us a comparison with Blount's result [8]. Another interesting feature of this approach is that it confirms the fundamental role played by Berry curvatures since it results in an effective diagonal Hamiltonian with Berry phase corrections as well

as noncommutative (Berry connections dependent) coordinates and momentum covariant operators as in previous approaches [5][6]. The resulting generic equations of motion are also corrected by Berry curvatures terms.

The paper is organized as follows. In the next section we develop the differential calculus in noncommutative space. We then derive, in section 3, the diagonalization procedure and show the link with the differential equation in section 4. Section 5 is devoted to the discussion of the dynamical coordinate and momentum operators. In section 6, we give the general diagonal energy operator formula to the second order in \hbar and in section 7 the case of the Dirac electron in an external electric field is considered as an illustration of the general formalism. Last section is for the conclusion.

II. PRELIMINARIES:

A. Differential calculus on noncommutative space

Consider a quantum mechanical system whose state space is a tensor product $L^2(\mathcal{R}^3) \otimes V$ with V some internal space. In other words, the Hamiltonian can be written as a matrix $H(\mathbf{P}, \mathbf{R})$ of size $\dim V$ whose elements are operators depending on a couple of canonical operators $[R^i, P^j] = i\hbar\delta_{ij}$, the archetype example being the Dirac Hamiltonian with $V = C^4$. Our goal is to derive the formal diagonal representation of this matrix valued quantum Hamiltonian. By diagonal Hamiltonian it is meant an effective in-band Hamiltonian which has a matrix representation with block-diagonal matrix elements associated with energy band subspaces. As in paper [5] the procedure we propose for the removal of the interband matrix elements needs the use of some unconventional mathematics compared to the usual formalism of quantum mechanics. Since the Planck constant is here considered as a variable, operators do not satisfy the usual rules of quantum mechanics. Therefore before directly embarking on the diagonalization procedure we first introduce the required mathematical tools. For that purpose we start with some definitions and notations.

1. Running coordinate and momentum operators

To begin with, we introduce a formal space of non commuting infinitesimal operators $dX_\alpha^i \equiv \{dR_\alpha^i, dP_\alpha^i\}$ $i = 1, 2, 3$ indexed by a continuous parameter α , that satisfy the

following infinitesimal Heisenberg algebra (with a reversed sign) :

$$[dR_\alpha^i, dP_{\alpha'}^j] = -i d\alpha \delta_{\alpha, \alpha'} \delta_{ij} \quad \text{and} \quad [dR_\alpha^i, dR_{\alpha'}^j] = [dP_\alpha^i, dP_{\alpha'}^j] = 0. \quad (1)$$

Then, we define a set of running coordinate and momentum operators by witting the following formal sums :

$$R_\alpha^i = R^i - \int_\alpha^{\hbar} dR_\lambda^i, \quad \text{and} \quad P_\alpha^i = P^i - \int_\alpha^{\hbar} dP_\lambda^i \quad (2)$$

with the choice of convention $dR_\alpha^i = R_\alpha^i - R_{\alpha-d\alpha}^i$ and $dP_\alpha^i = P_\alpha^i - P_{\alpha-d\alpha}^i$. This "downward" choice of the differential element notably implies the commutation rules $[R_\alpha^i, dP_\alpha^j] = [dR_\alpha^i, P_\alpha^j] = 0$ which turns out to be absolutely necessary to develop later on a differential calculus on this noncommutative space.

For $\alpha = \hbar$ we recover the usual canonical operators $R^i \equiv R_\hbar^i$ and $P^i \equiv P_\hbar^i$ which satisfy the canonical Heisenberg algebra $[R^i, P^j] = i\hbar \delta_{ij}$, whereas the running operators satisfy

$$[R_\alpha^i, P_\alpha^j] = i\alpha \delta_{ij} \quad \text{and} \quad [R_\alpha^i, R_\alpha^j] = [P_\alpha^i, P_\alpha^j] = 0. \quad (3)$$

Note that in this paper we will never consider the algebra of the operators R_α^i and $P_{\alpha'}^j$ for $\alpha \neq \alpha'$, which from the definition Eq. (2) is clearly not a Heisenberg one.

The variables $d\mathbf{X}_\alpha$ have to be understood as fictitious variables that make the link between quantized operators ($\alpha = \hbar$) and classical variables ($\alpha = 0$). As will be explicit later on, their role is to transport our quantum system to a formal classical one that can be in general easily diagonalized, and then back from the formal classical one to the required quantized system. By writing $d\mathbf{X}_\alpha = \sqrt{d\alpha/\hbar} \hat{\mathbf{X}}_\alpha$ with $\hat{\mathbf{X}}_\alpha$ a normalized canonical operator we see that the infinitesimal quantities $d\mathbf{X}_\alpha$ are actually of order $\sqrt{d\alpha}$. However, having this in mind, we will never use the $\hat{\mathbf{X}}_\alpha$ notation and always work with $d\mathbf{X}_\alpha$.

2. Differential algebra

For the sequel, we need to define the differential $dF(\mathbf{X}_\alpha, \alpha)$ of an arbitrary function $F(\mathbf{X}_\alpha, \alpha)$ where $X_\alpha^i \equiv \{R_\alpha^i, P_\alpha^i\}$. For this purpose, consider the operators $R^i - \int_\alpha^{\hbar} dR_\lambda^i$, and $P_\alpha^i = P^i - \int_\alpha^{\hbar} dP_\lambda^i$ as acting on a space $W = (V \otimes L^2(\mathcal{R}^3)) \otimes (\otimes_{\alpha < \hbar} L^2(\mathcal{R}^3)_\alpha)$ which is the tensor product of V and an infinite number of copies of $L^2(\mathcal{R}^3)$. The tensor product $V \otimes L^2(\mathcal{R}^3)$ and the space $L^2(\mathcal{R}^3)_\alpha$ refer respectively to the spaces on which the canonical operators (R^i, P^i) and the differential operators dX_α^i act.

Now to be consistent with our convention for the differential element dX_α^i , the differential of the function $F(\mathbf{X}_\alpha, \alpha)$ is defined in an unusual backward manner as $dF(\mathbf{X}_\alpha, \alpha) \equiv F(\mathbf{X}_\alpha, \alpha) - F(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha)$. This construction is essential as it permits the commutation between the non-differential and the differential elements dX_α^i . The differential will be written as a second order expansion plus some neglected terms:

$$\begin{aligned} dF(\mathbf{X}_\alpha, \alpha) &= \nabla_{R_i} F dR_\alpha^i + \nabla_{P_i} F dP_\alpha^i - \frac{1}{2} (\nabla_{R_i} \nabla_{R_j} F) dR_\alpha^i dR_\alpha^j - \frac{1}{2} (\nabla_{P_i} \nabla_{P_j} F) dP_\alpha^i dP_\alpha^j \\ &\quad - \frac{1}{2} (\nabla_{R_i} \nabla_{P^i} F) (dR_\alpha^i dP_\alpha^j + dP_\alpha^j dR_\alpha^i) + \langle F(\mathbf{X}_\alpha, \alpha) \rangle d\alpha + \frac{\partial F}{\partial \alpha} d\alpha \\ &\quad + \text{terms of order 3,} \end{aligned} \tag{4}$$

where all expressions in the r.h.s. are evaluated at $(\mathbf{X}_\alpha, \alpha)$. Note that here, we have kept the terms of order square in dX_α^i since they are of order $d\alpha$ and thus contribute to the differential, whereas higher orders can safely be disregarded as they are negligible when the integration over α is considered. These second order terms have been organized in a certain form that will be practical for us later. There is nothing to say about the usual $dR_\alpha^i dR_\alpha^j$ or $dP_\alpha^i dP_\alpha^j$ terms. The crossed terms involving products of dP_α^i and dR_α^j have to be taken with care since these two infinitesimal elements do not commute with each other. For this purpose we have decomposed the second order terms involving products such as $dP_\alpha^i dR_\alpha^j$ and $dR_\alpha^j dP_\alpha^i$ in a symmetric part proportional to $(dR_\alpha^i dP_\alpha^j + dP_\alpha^j dR_\alpha^i)$ giving the contribution $-\frac{1}{2} (\nabla_{R_i} \nabla_{P^i} F) (dR_\alpha^i dP_\alpha^j + dP_\alpha^j dR_\alpha^i)$, and an antisymmetric part proportional to $(dR_\alpha^i dP_\alpha^j - dP_\alpha^j dR_\alpha^i) = -i\delta^{ij} d\alpha$ corresponding to the bracket $\langle F(\mathbf{X}_\alpha, \alpha) \rangle d\alpha$ that we now explain in details. The notation $\langle F(\mathbf{X}_\alpha, \alpha) \rangle$ (which in [5] was corresponding to the operation $-\frac{i}{2} \text{Asym} \nabla_{R_i} \nabla_{P^i} F(\mathbf{X}_\alpha, \alpha)$) is defined as a specific procedure on a series expansion of F in the variables R_α^i, P_α^i in the following way : let F be a sum of monomials of the kind $M_1(\mathbf{R}_\alpha) M_2(\mathbf{P}_\alpha) M_3(\mathbf{R}_\alpha) \dots$ the M_i being arbitrary monomials in R_α or P_α alternatively. Let the operator $\nabla_{R_i} \nabla_{P^i}$ acts on such an expression by deriving all combinations of one monomial in \mathbf{R}_α and one monomial in \mathbf{P}_α . For each of these combinations, insert a dR_α^i at the place where the derivative ∇_{R_i} is acting and in a same manner a dP_α^j at the place where the derivative ∇_{P^i} is acting. This leads to an expression with two kind of terms, one kind being proportional to the $dR_\alpha^i dP_\alpha^j$, and the second proportional to $dP_\alpha^j dR_\alpha^i$. Then rewrite this expression in terms of $dR_\alpha^i dP_\alpha^j + dP_\alpha^j dR_\alpha^i$ and $dR_\alpha^i dP_\alpha^j - dP_\alpha^j dR_\alpha^i = -i\delta^{ij} d\alpha$. Then $\langle F(\mathbf{X}_\alpha, \alpha) \rangle$ is defined as minus the contributions of terms proportional to $-i\delta^{ij} d\alpha$ in the computation in the procedure just considered. This definition implies a procedure which is

clearly dependent of the symmetrization chosen for the expansion of F .

To make the definition of $\langle F(\mathbf{X}_\alpha, \alpha) \rangle$ clearer, consider some important practical examples. If the function F has the following form $F = \frac{1}{2} (A(\mathbf{R}_\alpha) B(\mathbf{P}_\alpha) + B(\mathbf{P}_\alpha) A(\mathbf{R}_\alpha))$ which corresponds to a frequent choice of symmetrization in R_α and P_α , then $\langle F(\mathbf{X}_\alpha, \alpha) \rangle = \frac{i}{4} [A(\mathbf{R}_\alpha), B(\mathbf{P}_\alpha)]$. Another choice of symmetrization leads in general to a different result.

For instance, if we rewrite the same function F in a fully symmetrized form in R_α and P_α (that is invariant by all permutations in R_α and P_α) which is also often used, we have now have a different result since $\langle F(\mathbf{X}_\alpha, \alpha) \rangle = 0$.

Nevertheless, this dependence of $\langle F(\mathbf{X}_\alpha, \alpha) \rangle$ in the symmetrization choice is not astonishing at all. Actually changing the symmetrization of a function $F(\mathbf{X}_\alpha, \alpha)$ introduces some explicit terms in α which changes also the term $\partial_\alpha F d\alpha$ present in the differential Eq. (4). As a consequence, neither the partial derivative with respect to α , nor the bracket are invariant by a change of form. But, what is invariant is the sum $\partial_\alpha F + \langle F \rangle$.

We show this assertion by first giving a useful formula for the differential of a product of functions :

$$\begin{aligned}
d(F(\mathbf{X}_\alpha) G(\mathbf{X}_\alpha)) &= d(F(\mathbf{X}_\alpha)) G(\mathbf{X}_\alpha) + F(\mathbf{X}_\alpha) d(G(\mathbf{X}_\alpha)) \\
&\quad - \nabla_{R_i} F \nabla_{R_j} G dR_\alpha^i dR_\alpha^j - \nabla_{P_i} F \nabla_{P_j} G dP_\alpha^i dP_\alpha^j \\
&\quad - (\nabla_{R_i} F \nabla_{P_j} G) dR_\alpha^i dP_\alpha^j - (\nabla_{P_j} F \nabla_{R_i} G) dP_\alpha^j dR_\alpha^i \\
&= d(F(\mathbf{X}_\alpha)) G(\mathbf{X}_\alpha) + F(\mathbf{X}_\alpha) d(G(\mathbf{X}_\alpha)) \\
&\quad - \nabla_{R_i} F \nabla_{R_j} G dR_\alpha^i dR_\alpha^j - \frac{1}{2} \nabla_{P_i} F \nabla_{P_j} G dP_\alpha^i dP_\alpha^j \\
&\quad - \frac{1}{2} (\nabla_{R_i} F \nabla_{P_j} G + \nabla_{P_j} F \nabla_{R_i} G) (dR_\alpha^i dP_\alpha^j + dP_\alpha^j dR_\alpha^i) \\
&\quad - \frac{i}{2} (\nabla_{P_i} F(\mathbf{X}_\alpha) \nabla_{R_i} G(\mathbf{X}_\alpha) - \nabla_{R_i} F(\mathbf{X}_\alpha) \nabla_{P_i} G(\mathbf{X}_\alpha)) d\alpha \quad (5)
\end{aligned}$$

Now, the invariance of $\partial_\alpha F + \langle F \rangle$ is the consequence of two facts. The first one is that the definition of the differential $dF(\mathbf{X}_\alpha, \alpha)$ is independent of the way the function $F(\mathbf{X}_\alpha, \alpha)$ has been symmetrized. Actually, considering $F(\mathbf{X}_\alpha, \alpha)$ as a monomial in the canonical variables, a change of symmetrization amounts to move successively powers of the momentum on the right or the left of the position variable. Recursively, one just has to check that moving only one power of the momentum does not change the differential. Considering thus a monomial of the form $F(\mathbf{X}_\alpha, \alpha) = M_1 P_\alpha^i R_\alpha^i M_2$ with M_1 and M_2 arbitrary, move the momentum to rewrite $F(\mathbf{X}_\alpha, \alpha) = M_1 (R_\alpha^i P_\alpha^i - i\alpha) M_2$. Our formula for a product of differential, allows to

assert that the differentials of the two terms $M_1 P_\alpha^i R_\alpha^i M_2$ and $M_1 (R_\alpha^i P_\alpha^i - i\alpha) M_2$ differ only by the contributions $M_1 d(P_\alpha^i R_\alpha^i) M_2$ and $M_1 d(R_\alpha^i P_\alpha^i - i\alpha) M_2$. As a consequence, the differential $dF(\mathbf{X}_\alpha, \alpha)$ is independent of the choice of symmetrization if and only if $d(P_\alpha^i R_\alpha^i) = d(R_\alpha^i P_\alpha^i - i\alpha)$. Given that

$$\begin{aligned} d(P_\alpha^i R_\alpha^i) &= dP_\alpha^i R_\alpha^i + P_\alpha^i dR_\alpha^i - \frac{1}{2} (dR_\alpha^i dP_\alpha^j + dP_\alpha^j dR_\alpha^i) - \frac{i}{2} d\alpha \\ d(R_\alpha^i P_\alpha^i - i\alpha) &= R_\alpha^i dP_\alpha^i + dR_\alpha^i P_\alpha^i - \frac{1}{2} (dR_\alpha^i dP_\alpha^j + dP_\alpha^j dR_\alpha^i) + \frac{i}{2} d\alpha - i d\alpha \end{aligned}$$

and that the dP_α^i, dR_α^i commute with R_α^i, P_α^i the equality of the two previous lines holds as well as the assertion of symmetrization independence.

This result joined to the second fact that in the series expansion of $dF(\mathbf{X}_\alpha, \alpha)$, the variables $dX_\alpha^i, dX_\alpha^i dX_\alpha^j, d\alpha$, are independent, and thus that the coefficients of the expansion of $dF(\mathbf{X}_\alpha, \alpha)$ are uniquely defined, leads directly to the announced result that $\partial_\alpha F + \langle F \rangle$ which is the coefficient of $d\alpha$ in the expansion of $dF(\mathbf{X}_\alpha, \alpha)$ is independent of any symmetrization choice. Therefore it is convenient to introduce a symmetrization invariant derivative D_α as given by

$$D_\alpha F(\mathbf{X}_\alpha, \alpha) = \partial_\alpha F(\mathbf{X}_\alpha, \alpha) + \langle F(\mathbf{X}_\alpha, \alpha) \rangle \quad (6)$$

We end up this paragraph by introducing two formulas that will be convenient later. First, we still more compactify our notations for the differential by writing :

$$\begin{aligned} dF(\mathbf{X}_\alpha, \alpha) &= \sum_{i=1}^6 \nabla_{X_\alpha^i} F(\mathbf{X}_\alpha, \alpha) dX_\alpha^i - \frac{1}{4} \sum_{i,j=1}^6 \nabla_{X_\alpha^i} \nabla_{X_\alpha^j} F(\mathbf{X}_\alpha, \alpha) (dX_\alpha^i dX_\alpha^j + dX_\alpha^j dX_\alpha^i) \\ &\quad + D_\alpha F(\mathbf{X}_\alpha, \alpha) d\alpha \end{aligned} \quad (7)$$

with $i, j = 1..6$. We also assume that $X_\alpha^i \equiv R_\alpha^i$ for $i = 1, 2, 3$ and $X_\alpha^i \equiv P_\alpha^i$ for $i = 4, 5, 6$. In fact, as we will show in the next sections the quantity of real importance for us in Eq. (11) is the term proportional to $d\alpha$. Note that in this notation, our previous formula for a product of differential takes a compact form :

$$\begin{aligned} d(F(\mathbf{X}_\alpha) G(\mathbf{X}_\alpha)) &= d(F(\mathbf{X}_\alpha)) G(\mathbf{X}_\alpha) + F(\mathbf{X}_\alpha) d(G(\mathbf{X}_\alpha)) \\ &\quad - \sum_{i,j=1}^6 \nabla_{X_\alpha^i} F(\mathbf{X}_\alpha, \alpha) \nabla_{X_\alpha^j} G(\mathbf{X}_\alpha, \alpha) dX_\alpha^i dX_\alpha^j \end{aligned}$$

where the last term could of course be developed as before in symmetric and antisymmetric parts.

Second, we also give the bracket formula $\langle \cdot \rangle$ for a product of two functions. Using the procedure defined previously (see also ref. [5]) one obtains the following expression

$$\begin{aligned} \langle F(\mathbf{X}_\alpha) G(\mathbf{X}_\alpha) \rangle &= \langle F(\mathbf{X}_\alpha) \rangle G(\mathbf{X}_\alpha) + F(\mathbf{X}_\alpha) \langle G(\mathbf{X}_\alpha) \rangle - \frac{i}{2} (\nabla_{P_i} F(\mathbf{X}_\alpha) \nabla_{R_i} G(\mathbf{X}_\alpha) \\ &\quad - \nabla_{R_i} F(\mathbf{X}_\alpha) \nabla_{P_i} G(\mathbf{X}_\alpha)). \end{aligned} \quad (8)$$

This formula shows that the bracket operation for a product $F(\mathbf{X}_\alpha) G(\mathbf{X}_\alpha)$ can also be seen as a sort of deformation of the Poisson bracket including some "internal" contributions $\langle F(\mathbf{X}_\alpha) \rangle$ and $\langle G(\mathbf{X}_\alpha) \rangle$.

Let us remark ultimately that the term $\langle F(\mathbf{X}_\alpha, \alpha) \rangle d\alpha$ in Eq. (4) which is of the second order in the derivatives ∇_{R_i} and ∇_{P_i} is very reminiscent of the bracket introduced in the stochastic calculus. Given the non commutativity of the operators at stake, one should in fact rather expect our formalism to be close to the quantum stochastic calculus [14]. However, it does not seem at first sight that our objects fit in such a framework which deals rather with a formalism of creation, annihilation and conservation operators. A full comparison is out of the scope of our paper.

3. Integration

Now, we will prove that the formula Eq. (4) allows to express a function $F(\mathbf{X}_\hbar, \hbar)$ which depends on the physical quantum operators X_\hbar defined at the quantum scale $\alpha = \hbar$, as the integration of the differential:

$$F(\mathbf{X}_\hbar, \hbar) = F(\mathbf{X}_0, 0) + \int_0^\hbar dF(\mathbf{X}_\alpha, \alpha) \quad (9)$$

where \mathbf{X}_0 are therefore to be considered as classical commuting quantities.

The proof of the previous formula is as before recursive. Expanding $F(\mathbf{X}_\alpha, \alpha)$ as sum of monomials in the canonical variable, it is enough to prove our assertion for a monomial of a certain degree in the R_\hbar^i, P_\hbar^i . For monomials of degree one, that is linear expressions in the R_\hbar^i, P_\hbar^i , the result is trivial given our definitions. Now, assume that the result is true for all monomial M of, say, bidegree m, n in the R_\hbar^i, P_\hbar^i . We will show the result for a monomial whose degree has increased by one in one variable. Such a monomial can be written $X_\hbar^i M$ or $M X_\hbar^i$, with M of bidegree m, n (recall that $X_\hbar^i = (R_\hbar^i, P_\hbar^i)$). We will concentrate on the

first possibility, the proof being obviously the same for the other case. We compute directly $X_0^i M(\mathbf{X}_0, 0) + \int_0^{\hbar} dX_\alpha^i M(\mathbf{X}_\alpha, \alpha)$.

Given our formula for the differential of a product as well as the recurrence hypothesis which states that the integral formula is true for M we have :

$$\begin{aligned} X_0^i M(\mathbf{X}_0, 0) + \int_0^{\hbar} d(X_\alpha^i M(\mathbf{X}_\alpha, \alpha)) &= \left(X_{\hbar}^i - \int_0^{\hbar} dX_\alpha^i \right) \left(M(X_{\hbar}^i, \hbar) - \int_0^{\hbar} dM(\mathbf{X}_\alpha, \alpha) \right) \\ &+ \int_0^{\hbar} dX_\alpha^i M(\mathbf{X}_\alpha, \alpha) + \int_0^{\hbar} X_\alpha^i dM(\mathbf{X}_\alpha, \alpha) \\ &- \sum_j \int_0^{\hbar} dX_\alpha^i dX_\alpha^j \nabla_{X_\alpha^j} M(\mathbf{X}_\alpha, \alpha) \end{aligned}$$

In the last line we have chosen (for the sake of simplicity) not to separate the product $dX_\alpha^i dX_\alpha^j$ into symmetric and antisymmetric part, but to keep the second order terms in a compact form. Note again the minus sign arising in front of this term due to our choice of definition for the differential. As a consequence one has :

$$\begin{aligned} X_0^i M(\mathbf{X}_0, 0) + \int_0^{\hbar} d(X_\alpha^i M(\mathbf{X}_\alpha, \alpha)) \\ = X_{\hbar}^i M(X_{\hbar}^i, \hbar) - \int_0^{\hbar} dX_\alpha^i M(X_{\hbar}^i, \hbar) - X_{\hbar}^i \int_0^{\hbar} dM(\mathbf{X}_\alpha, \alpha) + \int_0^{\hbar} dX_\alpha^i \int_0^{\hbar} dM(\mathbf{X}_\alpha, \alpha) \\ + \int_0^{\hbar} dX_\alpha^i M(\mathbf{X}_\alpha, \alpha) + \int_0^{\hbar} X_\alpha^i dM(\mathbf{X}_\alpha, \alpha) - \sum_j \int_0^{\hbar} dX_\alpha^i dX_\alpha^j \nabla_{X_\alpha^j} M(\mathbf{X}_\alpha, \alpha) \end{aligned}$$

Now, due again to the recursion hypothesis, we rewrite $\int_0^{\hbar} dX_\alpha^i M(\mathbf{X}_\alpha, \alpha) + \int_0^{\hbar} X_\alpha^i dM(\mathbf{X}_\alpha, \alpha)$ in the following manner :

$$\begin{aligned} \int_0^{\hbar} dX_\alpha^i M(\mathbf{X}_\alpha, \alpha) + \int_0^{\hbar} X_\alpha^i dM(\mathbf{X}_\alpha, \alpha) &= \int_0^{\hbar} dX_\alpha^i \left[M(X_{\hbar}^i, \hbar) - \int_\alpha^{\hbar} dM(\mathbf{X}_{\alpha'}, \alpha') \right] \\ &+ \int_0^{\hbar} \left(X_{\hbar}^i - \int_\alpha^{\hbar} dX_{\alpha'}^i \right) dM(\mathbf{X}_\alpha, \alpha) \\ &= \int_0^{\hbar} dX_\alpha^i M(X_{\hbar}^i, \hbar) + X_{\hbar}^i \int_0^{\hbar} dM(\mathbf{X}_\alpha, \alpha) \\ &- \int_0^{\hbar} dX_\alpha^i \left[\int_\alpha^{\hbar} dM(\mathbf{X}_{\alpha'}, \alpha') \right] - \int_0^{\hbar} \int_\alpha^{\hbar} dX_{\alpha'}^i dM(\mathbf{X}_\alpha, \alpha) \end{aligned}$$

Now, we have to take care about the meaning of the two last integrals. Actually, given our definition of the differential, in the first integral $\int_0^{\hbar} dX_\alpha^i \left[\int_\alpha^{\hbar} dM(\mathbf{X}_{\alpha'}, \alpha') \right]$, one has to consider that $\alpha' > \alpha$, since our differentials are pointing downward, and as a consequence when α' is getting closer to α , the "last" element of integration is $dM(\mathbf{X}_{\alpha+d\alpha}, \alpha+d\alpha) =$

$M(\mathbf{X}_{\alpha+d\alpha}, \alpha + d\alpha) - M(\mathbf{X}_\alpha, \alpha)$. In the integral $\int_0^{\hbar} \int_\alpha^{\hbar} dX_\alpha^i dM(\mathbf{X}_\alpha, \alpha)$ one has rather to consider $\alpha' < \alpha$. As a consequence, one has :

$$\begin{aligned} & \int \int_{\alpha' < \alpha < \hbar} dX_\alpha^i dM(\mathbf{X}_{\alpha'}, \alpha') + \int \int_{\alpha < \alpha' < \hbar} dX_\alpha^i dM(\mathbf{X}_{\alpha'}, \alpha') \\ &= \left(\int_0^{\hbar} dX_\alpha^i \int_0^{\hbar} dM(\mathbf{X}_\alpha, \alpha) - \int_0^{\hbar} dX_\alpha^i dM(\mathbf{X}_\alpha, \alpha) \right) \end{aligned}$$

the last term in the right cancelling the diagonal contribution that should not appear given our considerations just above. Note that this contribution would be negligible for an ordinary integral. As a consequence, we obtain :

$$\begin{aligned} \int_0^{\hbar} dX_\alpha^i M(\mathbf{X}_\alpha, \alpha) + \int_0^{\hbar} X_\alpha^i dM(\mathbf{X}_\alpha, \alpha) &= \int_0^{\hbar} dX_\alpha^i M(X_\hbar^i, \hbar) + X_\hbar^i \int_0^{\hbar} dM(\mathbf{X}_\alpha, \alpha) \\ &\quad - \iint_{\alpha' < \alpha < \hbar} dX_\alpha^i dM(\mathbf{X}_{\alpha'}, \alpha') - \iint_{\alpha < \alpha' < \hbar} dX_\alpha^i dM(\mathbf{X}_{\alpha'}, \alpha') \\ &= \int_0^{\hbar} dX_\alpha^i M(X_\hbar^i, \hbar) + X_\hbar^i \int_0^{\hbar} dM(\mathbf{X}_\alpha, \alpha) \\ &\quad - \left(\int_0^{\hbar} dX_\alpha^i \int_0^{\hbar} dM(\mathbf{X}_\alpha, \alpha) - \int_0^{\hbar} dX_\alpha^i dM(\mathbf{X}_\alpha, \alpha) \right) \end{aligned}$$

Gathering all the terms leads thus ultimately to :

$$\begin{aligned} X_0^i M(\mathbf{X}_0, 0) + \int_0^{\hbar} d(X_\alpha^i M(\mathbf{X}_\alpha, \alpha)) &= X_\hbar^i M(X_\hbar^i, \hbar) \\ &\quad + \int_0^{\hbar} dX_\alpha^i dM(\mathbf{X}_\alpha, \alpha) - \sum_j \int_0^{\hbar} dX_\alpha^i dX_\alpha^j \nabla_{X_\alpha^j} M(\mathbf{X}_\alpha, \alpha) \end{aligned}$$

Now, the integral on the diagonal : $\int_0^{\hbar} dX_\alpha^i dM(\mathbf{X}_\alpha, \alpha)$ reduces to $\sum_j \int_0^{\hbar} dX_\alpha^i dX_\alpha^j \nabla_{X_\alpha^j} M(\mathbf{X}_\alpha, \alpha)$. Actually as explained before in the definition of the differential, the second order terms in $dM(\mathbf{X}_\alpha, \alpha)$, multiplied by dX_α^i contribute to the order $d\alpha^{\frac{3}{2}}$ which yields a zero contribution while summing over α . We thus end up with :

$$X_0^i M(\mathbf{X}_0, 0) + \int_0^{\hbar} d(X_\alpha^i M(\mathbf{X}_\alpha, \alpha)) = X_\hbar^i M(X_\hbar^i, \hbar) \quad (10)$$

which proves Eq. (9).

Now, we interpret Eq. (9) as follows. Starting from the classical regime the integration over α leads in continuous manner to the fully quantum regime. This original manner to achieve a sort of quantification is very reminiscent of the renormalization group technics in

which the integration of the high energy modes leads to an effective low energy theory with all quantum fluctuations included.

Eq. (9) allows us to rewrite formally a function $F(\mathbf{X}_{\hbar}, \hbar)$ depending on the physical variables in terms of the same function evaluated at the classical variables $(\mathbf{X}_0, 0)$ plus an integral of a differential. This introduces apparently a useless complexity since the expansion of the right hand side of Eq. (9) around \mathbf{X}_{\hbar} involves an infinite series of the infinitesimal increment of the $d\mathbf{X}_{\alpha}$. By construction, this expansion implies trivially that all terms of the series exactly vanish except obviously the initial term which is equal to the function $F(\mathbf{X}_{\hbar}, \hbar)$. Although trivial, this decomposition will prove useful later.

Note again that the integral in Eq. (9) has to be understood as being computed downward as seen in the definition of $(\mathbf{R}_{\alpha}, \mathbf{P}_{\alpha})$: the starting point is at \hbar and the differentials are pointed downward. For example $\nabla_{\mathbf{R}_{\alpha}} F \cdot d\mathbf{R}_{\alpha} = -\nabla_{\mathbf{R}_{\alpha}} F(\mathbf{X}_{\alpha}, \alpha) (\mathbf{R}_{\alpha-d\alpha} - \mathbf{R}_{\alpha})$. However, this equation has to be taken with some care. Actually the sum over the terms proportional to $d\mathbf{X}_{\alpha}$ being a sum of terms of magnitude $\sqrt{d\alpha}$, it converges only if the sum is discretized, and this will be implicitly assumed in this paper. Defining properly the continuous limit is out of the scope of this paper.

For later convenience we still more compactify our notations by writing :

$$dF(\mathbf{X}_{\alpha}, \alpha) = \sum_{i=1}^6 \nabla_{X_{\alpha}^i} F(\mathbf{X}_{\alpha}, \alpha) dX_{\alpha}^i - \frac{1}{4} \sum_{i,j=1}^6 \nabla_{X_{\alpha}^i} \nabla_{X_{\alpha}^j} F(\mathbf{X}_{\alpha}, \alpha) (dX_{\alpha}^i dX_{\alpha}^j + dX_{\alpha}^j dX_{\alpha}^i) + D_{\alpha} F(\mathbf{X}_{\alpha}, \alpha) d\alpha \quad (11)$$

with $i, j = 1..6$. We also assume that $X_{\alpha}^i \equiv R_{\alpha}^i$ for $i = 1, 2, 3$ and $X_{\alpha}^i \equiv P_{\alpha}^i$ for $i = 4, 5, 6$. In fact, as we will show in the next sections the quantity of real importance for us in Eq. (11) is the term proportional to $d\alpha$.

For later use we now give the bracket formula $\langle \cdot \rangle$ for a product of two functions. Using the procedure defined previously (see also ref. [5]) one obtains the following expression

$$\langle F(\mathbf{X}_{\alpha}) G(\mathbf{X}_{\alpha}) \rangle = \langle F(\mathbf{X}_{\alpha}) \rangle G(\mathbf{X}_{\alpha}) + F(\mathbf{X}_{\alpha}) \langle G(\mathbf{X}_{\alpha}) \rangle - \frac{i}{2} (\nabla_{P_i} F(\mathbf{X}_{\alpha}) \nabla_{R_i} G(\mathbf{X}_{\alpha}) - \nabla_{R_i} F(\mathbf{X}_{\alpha}) \nabla_{P_i} G(\mathbf{X}_{\alpha})). \quad (12)$$

This formula shows that the bracket operation for a product $F(\mathbf{X}_{\alpha}) G(\mathbf{X}_{\alpha})$ can also be seen as a sort of deformation of the Poisson bracket including some "internal" contributions $\langle F(\mathbf{X}_{\alpha}) \rangle$ and $\langle G(\mathbf{X}_{\alpha}) \rangle$.

Before going to the diagonalization procedure we still need to define an expectation operator which will allow us to express in a specific way the same function F evaluated at $(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha)$ and $(\mathbf{X}_\alpha, \alpha - d\alpha)$.

B. Expectation operator

Having now a differential set up, we aim at defining a linear conditional expectation operator $\mathcal{E}(\cdot)$. We first set the formula :

$$\mathcal{E}(d\alpha) = d\alpha \quad \text{and} \quad \mathcal{E}(F(\mathbf{X}_{\hbar}, \hbar)) = F(\mathbf{X}_{\hbar}, \hbar) \quad (13)$$

The first equality will allow to define expectations of integrals of a function $F(\mathbf{X}_\alpha, \alpha)$ with a deterministic increment $d\alpha$. The second one will make the expectation conditional by fixing the starting point \mathbf{X}_{\hbar} of the path \mathbf{X}_α . By analogy with the stochastic calculus, \mathbf{X}_α will be seen as a random path whose infinitesimal increments $d\mathbf{X}_\alpha$ will have zero expectation. However, again as in stochastic calculus, the expectation of a quadratic term like $dX_\alpha^i dX_\alpha^j$ can not be set consistently to vanish since it can be written in terms of $\mathcal{E}(dR_\alpha^i dP_\alpha^j + dP_\alpha^j dR_\alpha^i)$ and $\mathcal{E}(dR_\alpha^i dP_\alpha^j - dP_\alpha^j dR_\alpha^i) = -i\delta^{ij}d\alpha$, and the last term can not vanish due to the first definition in Eq.(13), and only the symmetric part $\mathcal{E}(dR_\alpha^i dP_\alpha^j + dP_\alpha^j dR_\alpha^i)$ can be chosen to cancel.

Consequently we set the expectation rules for products at the same "time" α

$$\mathcal{E}\left(\prod_{i=1}^n dR_{\alpha_i}\right) = \mathcal{E}\left(\prod_{i=1}^n dP_{\alpha_i}\right) = 0, \quad \mathcal{E}\left(\prod_{i=1}^n (dX_{\alpha_i}^{k_i} dX_{\alpha_i}^{l_i} + dX_{\alpha_i}^{l_i} dX_{\alpha_i}^{k_i})\right) = 0 \quad (14)$$

and for different times

$$\mathcal{E}\left(\prod_{i=1}^n \prod_{j=1}^p dX_{\alpha_i}^{k_i} dX_{\alpha_j}^{l_j}\right) = 0 \quad \text{for} \quad \alpha_i \neq \alpha_j \quad (15)$$

We need also the independence of the increment $d\mathbf{X}_\alpha$ with respect to a function evaluated at the corresponding \mathbf{X}_α which is formulated as

$$\mathcal{E}(F(\mathbf{X}_\alpha) d\mathbf{X}_\alpha) = \mathcal{E}(F(\mathbf{X}_\alpha)) \mathcal{E}(d\mathbf{X}_\alpha) = 0$$

A first consequence of these definitions is

$$\begin{aligned} \mathcal{E}(dF(\mathbf{X}_\alpha, \alpha)) &= \mathcal{E}(\partial_\alpha F(\mathbf{X}_\alpha, \alpha) + \langle F(\mathbf{X}_\alpha, \alpha) \rangle) d\alpha \\ &= \mathcal{E}(D_\alpha(\mathbf{X}_\alpha, \alpha)) d\alpha \end{aligned}$$

The main interest of these definitions will show up later on when a function $F(\mathbf{X}_{\hbar}, \hbar)$ will be written as an series expansion of iterated integrals along the path \mathbf{X}_{α} . Indeed, as already said before, the expansion of the two terms in the right hand side of Eq. (9) around \mathbf{X}_{\hbar} , implies a mechanism of cancellation of all contributions proportional to products of $dX_{\alpha_i}^{k_i}$. The expectation operator is built to cancel separately the contributions of products of $dX_{\alpha_i}^{k_i}$ at different times and symmetric products ($dX_{\alpha_i}^{k_i}dX_{\alpha_i}^{l_i} + dX_{\alpha_i}^{l_i}dX_{\alpha_i}^{k_i}$) at the same time, coming from the two terms of the right hand side of Eq. (9). The rest of the contributions form the right hand of Eq. (9) proportional to $d\alpha$ cancel exactly. As a consequence of this, an arbitrary function can be written

$$\begin{aligned} F_{\hbar}(\mathbf{X}_{\hbar}) &= \mathcal{E}(F_{\hbar}(\mathbf{X}_{\hbar})) = \mathcal{E}\left(F(\mathbf{X}_0, 0) + \int_0^{\hbar} dF(\mathbf{X}_{\alpha}, \alpha)\right) \\ &= \mathcal{E}(F(\mathbf{X}_0, 0)) + \mathcal{E}\left(\int_0^{\hbar} D_{\alpha}(\mathbf{X}_{\alpha}, \alpha) d\alpha\right) \end{aligned}$$

Once again we have apparently introduced some useless complexity, since the function $F_{\hbar}(\mathbf{X}_{\hbar})$ has been replaced by $\mathcal{E}(F(\mathbf{X}_0, 0))$ which cannot be computed directly, plus the integral $\mathcal{E}\left(\int_0^{\hbar} D_{\alpha}(\mathbf{X}_{\alpha}, \alpha) d\alpha\right)$. Irrelevant sums of infinitesimal increments have been introduced in both terms and thus ultimately have to cancel. However, this formulation presents the advantage of replacing \mathbf{X}_{\hbar} by \mathbf{X}_0 in the first contribution which will simplify some of our computations later. The integration over $d\alpha$ will allow to start a recursive expansion in the Planck constant. But to do so, we first need to introduce an other operator connecting a function evaluated at \mathbf{X}_0 to the same function evaluated at \mathbf{X}_{\hbar} . This is the role of the connection operator. We introduce it by relating two expectations of the same function evaluated at two different variables namely $\mathcal{E}(F(\mathbf{X}_{\alpha}, \alpha))$ to $\mathcal{E}(F(\mathbf{X}_{\alpha'}, \alpha))$ but with constant α . Given our previous definitions Eq. (11), one has

$$\begin{aligned} F(\mathbf{X}_{\alpha-d\alpha}, \alpha-d\alpha) - F(\mathbf{X}_{\alpha}, \alpha-d\alpha) &= -\langle F(\mathbf{X}_{\alpha}, \alpha-d\alpha) \rangle d\alpha - \sum_{i=1}^6 \nabla_{X_{\alpha}^i} F(\mathbf{X}_{\alpha}, \alpha-d\alpha) dX_{\alpha}^i \\ &+ \frac{1}{4} \sum_{i,j=1}^6 \nabla_{X_{\alpha}^j} \nabla_{X_{\alpha}^i} F(\mathbf{X}_{\alpha}, \alpha-d\alpha) (dX_{\alpha}^i dX_{\alpha}^j + dX_{\alpha}^j dX_{\alpha}^i) \end{aligned} \quad (16)$$

The absence of the $\partial/\partial\alpha$ term implies that this expression really depends on the choice of the symmetrization of the function $F(\mathbf{X}_{\alpha}, \alpha)$. As a consequence considering the expectation,

one has :

$$\mathcal{E} (F(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha)) = \mathcal{E} ((1 - \langle \cdot \rangle) d\alpha) F(\mathbf{X}_\alpha, \alpha - d\alpha) \quad (17)$$

where $\langle \cdot \rangle$ is the bracket operator defined previously. At this point we have to remark that, since the $\langle \cdot \rangle$ operation depends on the symmetrization procedure the expression on the right hand side of Eq.(17) depends on the choice of symmetrization in the variables \mathbf{X}_α in $F(\mathbf{X}_\alpha, \alpha)$. However, and this is the most important for us, the operation that sends $F(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha)$ to $(1 - \langle \cdot \rangle) d\alpha) F(\mathbf{X}_\alpha, \alpha - d\alpha)$ is independent at the first order in $d\alpha$ of the choice of symmetrization for the operators in the series expansion of $F(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha)$. This operation, that will be the only one relevant for us, is thus a consistent and independent of any symmetrization. The apparent trouble with the right hand side in Eq. (17) comes only from the fact that changing the order of operators $\mathbf{X}_{\alpha-d\alpha}$ is not the same operation as changing the order of operators \mathbf{X}_α due to different commutation relations. To avoid any confusion a subscript $\alpha - d\alpha$ should be added to F to remind the symmetrization has to be considered with variables $\mathbf{X}_{\alpha-d\alpha}$, but it would complicate the notations in a useless way for the present paper.

To check our last assertion of symmetrization independence, assume that $F(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha)$ is a monomial of the form $A(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha) R_{\alpha-d\alpha}^i P_{\alpha-d\alpha}^i B(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha)$ where A and B are arbitrary products of the $\mathbf{R}_{\alpha-d\alpha}$ and $\mathbf{P}_{\alpha-d\alpha}$. Assume now that we want to change the order of $R_{\alpha-d\alpha}^i$ and $P_{\alpha-d\alpha}^i$ in F , we just have to show that the operation considered $(1 - \langle \cdot \rangle) d\alpha) F(\mathbf{X}_\alpha, \alpha - d\alpha)$ is invariant under such exchange. Given our formula for the bracket of a product of several expressions Eq. (12), one has (we will omit the dependence in $\alpha - d\alpha$ in the right hand side since it plays no role in the argument) :

$$\begin{aligned} (1 - \langle \cdot \rangle) d\alpha) F(\mathbf{X}_\alpha, \alpha - d\alpha) &= A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i B(\mathbf{X}_\alpha) - \langle A(\mathbf{X}_\alpha) \rangle R_\alpha^i P_\alpha^i B(\mathbf{X}_\alpha) d\alpha \\ &\quad - A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \langle B(\mathbf{X}_\alpha) \rangle d\alpha - \frac{i}{2} A(\mathbf{X}_\alpha) d\alpha B(\mathbf{X}_\alpha) d\alpha \\ &\quad + \frac{i}{2} (\nabla_{P_i} A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \nabla_{R_i} B(\mathbf{X}_\alpha) - \nabla_{R_i} A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \nabla_{P_i} B(\mathbf{X}_\alpha)) d\alpha \\ &\quad + \frac{i}{2} (\nabla_{P_i} A(\mathbf{X}_\alpha) P_\alpha^i B(\mathbf{X}_\alpha) - \nabla_{R_i} A(\mathbf{X}_\alpha) R_\alpha^i B(\mathbf{X}_\alpha)) d\alpha \\ &\quad - \frac{i}{2} (A(\mathbf{X}_\alpha) P_\alpha^i \nabla_{R_i} B(\mathbf{X}_\alpha) - A(\mathbf{X}_\alpha) R_\alpha^i \nabla_{P_i} B(\mathbf{X}_\alpha)) d\alpha \quad (18) \end{aligned}$$

Note that all expressions in the right hand side are considered at the scale $\alpha - d\alpha$ ($A(\mathbf{X}_\alpha) \equiv A(\mathbf{X}_\alpha, \alpha - d\alpha)$ and idem for B).

On another hand, changing now the order of $R_{\alpha-d\alpha}^i$ and $P_{\alpha-d\alpha}^i$ yields

$$F(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha) = A(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha) P_{\alpha-d\alpha}^i R_{\alpha-d\alpha}^i B(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha) + i(\alpha - d\alpha) A(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha) B(\mathbf{X}_{\alpha-d\alpha}, \alpha - d\alpha) \quad (19)$$

For this particular symmetrization our operation $(1 - \langle \cdot \rangle d\alpha)$ gives now (skipping again the direct dependance in $\alpha - d\alpha$ that plays no role) :

$$\begin{aligned} (1 - \langle \cdot \rangle d\alpha) F(\mathbf{X}_\alpha, \alpha - d\alpha) &= A(\mathbf{X}_\alpha) P_\alpha^i R_\alpha^i B(\mathbf{X}_\alpha) + i(\alpha - d\alpha) A(\mathbf{X}_\alpha) B(\mathbf{X}_\alpha) \\ &+ \frac{i}{2} A(\mathbf{X}_\alpha) B(\mathbf{X}_\alpha) d\alpha - \langle A(\mathbf{X}_\alpha) \rangle P_\alpha^i R_\alpha^i B(\mathbf{X}_\alpha) d\alpha \\ &- A(\mathbf{X}_\alpha) P_\alpha^i R_\alpha^i \langle B(\mathbf{X}_\alpha) \rangle d\alpha - i(\alpha - d\alpha) \langle A(\mathbf{X}_\alpha) B(\mathbf{X}_\alpha) \rangle d\alpha \\ &+ \frac{i}{2} (\nabla_{P_i} A(\mathbf{X}_\alpha) P_\alpha^i R_\alpha^i \nabla_{R_i} B(\mathbf{X}_\alpha) - \nabla_{R_i} A(\mathbf{X}_\alpha) P_\alpha^i R_\alpha^i \nabla_{P_i} B(\mathbf{X}_\alpha)) d\alpha \\ &+ \frac{i}{2} (\nabla_{P_i} A(\mathbf{X}_\alpha) P_\alpha^i B(\mathbf{X}_\alpha) - \nabla_{R_i} A(\mathbf{X}_\alpha) R_\alpha^i B(\mathbf{X}_\alpha)) d\alpha \\ &- \frac{i}{2} (A(\mathbf{X}_\alpha) P_\alpha^i \nabla_{R_i} B(\mathbf{X}_\alpha) - A(\mathbf{X}_\alpha) R_\alpha^i \nabla_{P_i} B(\mathbf{X}_\alpha)) d\alpha \quad (20) \end{aligned}$$

We now rewrite the first three terms $Q_1 = A(\mathbf{X}_\alpha, \alpha) P_\alpha^i R_\alpha^i B(\mathbf{X}_{\alpha-d\alpha}, \alpha) + i(\alpha - d\alpha) A(\mathbf{X}_\alpha) B(\mathbf{X}_\alpha) + \frac{i}{2} A(\mathbf{X}_\alpha) B(\mathbf{X}_\alpha) d\alpha$ in the right hand side of Eq.(20) in the following manner $Q_1 = AR_\alpha^i P_\alpha^i B - id\alpha AB + \frac{i}{2} ABd\alpha$ which can also be written :

$$\begin{aligned} Q_1 &= A(\mathbf{X}_\alpha, \alpha) R_\alpha^i P_\alpha^i B(\mathbf{X}_{\alpha-d\alpha}, \alpha) - \frac{i}{2} A(\mathbf{X}_\alpha) d\alpha B(\mathbf{X}_\alpha) - \langle A(\mathbf{X}_\alpha) \rangle R_\alpha^i P_\alpha^i B(\mathbf{X}_\alpha) \\ &- A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \langle B(\mathbf{X}_\alpha) \rangle + \frac{i}{2} A(\mathbf{X}_\alpha) B(\mathbf{X}_\alpha) d\alpha \quad (21) \end{aligned}$$

In the same manner let rewrite the following contributions to Eq. (20) $Q_2 = -\langle A \rangle P_\alpha^i R_\alpha^i B d\alpha - AP_\alpha^i R_\alpha^i \langle B \rangle d\alpha - i(\alpha - d\alpha) \langle AB \rangle d\alpha + \frac{i}{2} (\nabla_{P_i} AP_\alpha^i R_\alpha^i \nabla_{R_i} B - \nabla_{R_i} AP_\alpha^i R_\alpha^i \nabla_{P_i} B)$ under the form

$$\begin{aligned} Q_2 &= -\langle A(\mathbf{X}_\alpha) \rangle R_\alpha^i P_\alpha^i B(\mathbf{X}_\alpha) d\alpha - A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \langle B(\mathbf{X}_\alpha) \rangle d\alpha - i(\alpha - d\alpha) \langle A(\mathbf{X}_\alpha) B(\mathbf{X}_\alpha, \alpha) \rangle d\alpha \\ &+ \frac{i}{2} (\nabla_{P_i} A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \nabla_{R_i} B(\mathbf{X}_\alpha) - \nabla_{R_i} A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \nabla_{P_i} B(\mathbf{X}_\alpha)) d\alpha \\ &+ \frac{\alpha}{2} (\nabla_{P_i} A(\mathbf{X}_\alpha) \nabla_{R_i} B(\mathbf{X}_\alpha) - \nabla_{R_i} A(\mathbf{X}_\alpha) \nabla_{P_i} B(\mathbf{X}_\alpha)) d\alpha \\ &+ i\alpha \langle A(\mathbf{X}_\alpha) \rangle B(\mathbf{X}_\alpha) + A(\mathbf{X}_\alpha) \langle B(\mathbf{X}_\alpha) \rangle d\alpha \end{aligned}$$

or

$$\begin{aligned}
Q_2 = & - \langle A(\mathbf{X}_\alpha) \rangle R_\alpha^i P_\alpha^i B(\mathbf{X}_\alpha) d\alpha - A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \langle B(\mathbf{X}_\alpha) \rangle d\alpha \\
& + \frac{i}{2} \left(\nabla_{P_i} A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \nabla_{R_i} B(\mathbf{X}_\alpha) - \nabla_{R_i} A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \nabla_{P_i} B(\mathbf{X}_\alpha) \right) d\alpha \\
& + i\alpha \langle A(\mathbf{X}_\alpha) B(\mathbf{X}_\alpha, \alpha) \rangle d\alpha - i(\alpha - d\alpha) \langle A(\mathbf{X}_\alpha) B(\mathbf{X}_\alpha, \alpha) \rangle d\alpha
\end{aligned}$$

At first order in $d\alpha$, Q_2 can then be written

$$\begin{aligned}
Q_2 = & - \langle A(\mathbf{X}_\alpha) \rangle R_\alpha^i P_\alpha^i B(\mathbf{X}_\alpha) d\alpha - A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \langle B(\mathbf{X}_\alpha) \rangle d\alpha \\
& + \frac{i}{2} \left(\nabla_{P_i} A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \nabla_{R_i} B(\mathbf{X}_\alpha) - \nabla_{R_i} A(\mathbf{X}_\alpha) R_\alpha^i P_\alpha^i \nabla_{P_i} B(\mathbf{X}_\alpha) \right) d\alpha \quad (22)
\end{aligned}$$

Gathering the various terms Q_1 and Q_2 yields the announced result of initial symmetrization independence as Eq. (20) is now the same as Eq. (18).

The meaning of our operator $1 - \langle \cdot \rangle d\alpha$ being now clarified, the above differential equation Eq. (17) can be easily integrated to get the required relation :

$$\mathcal{E}(F(\mathbf{X}_{\alpha_2}, \alpha_2)) = \mathcal{E} \left(\left[T \exp \left(- \int_{\alpha_2}^{\alpha_1} [S_{\mathbf{X}_{\alpha_1}} \langle \cdot \rangle_\alpha S_{\mathbf{X}_\alpha}] d\alpha \right) F(\mathbf{X}_{\alpha_2}, \alpha_2) \right] \right) \quad (23)$$

valid for $\alpha_1 > \alpha_2$ and the bracket $\langle \cdot \rangle_\alpha$ reminding that all commutators between canonical variables have to be computed with values $i\alpha$. The operator $S_{\mathbf{X}_\alpha}$ is a shift operator that sets the dynamical variables to \mathbf{X}_α . that is $S_{\mathbf{X}_\alpha} F(\mathbf{X}_\beta) = F(\mathbf{X}_\alpha)$. Note that we define it as acting backward or forward, i.e. $S_{\mathbf{X}_\alpha} S_{\mathbf{X}_\beta} = S_{\mathbf{X}_\alpha}$ whatever the values of α and β . In addition, T is the notation for the ordered exponential, that is :

$$\begin{aligned}
T \exp \left(- \int_{\alpha_2}^{\alpha_1} S_{\mathbf{X}_{\alpha_1}} \langle \cdot \rangle_\alpha S_{\mathbf{X}_\alpha} d\alpha \right) &= \sum_{\alpha_2 < \beta_n < \dots < \beta_1 < \alpha_1} \int \left[-S_{\mathbf{X}_{\alpha_1}} \langle \cdot \rangle_{\beta_n} S_{\mathbf{X}_{\beta_n}} \right] \dots \left[-S_{\mathbf{X}_{\alpha_1}} \langle \cdot \rangle_{\beta_1} S_{\mathbf{X}_{\beta_1}} \right] d\beta_1 \dots d\beta_n \\
&= \sum_{\alpha_2 < \beta_n < \dots < \beta_1 < \alpha_1} \int \left[-S_{\mathbf{X}_{\alpha_1}} \langle \cdot \rangle_{\beta_n} S_{\mathbf{X}_{\beta_n}} \right] \dots \left[-\langle \cdot \rangle_{\beta_1} S_{\mathbf{X}_{\beta_1}} \right] d\beta_1 \dots d\beta_n
\end{aligned}$$

Once again, the result is independent of the choice of initial symmetrization chosen in the variables \mathbf{X}_{α_2} . The introduction of the shift $S_{\mathbf{X}_\alpha}$ is crucial and induced by our previous considerations since the bracket operation has to be accompanied by a shift of the variable all along the integration process, so that ultimately \mathbf{X}_{α_2} is replaced by \mathbf{X}_{α_1} .

To gain some space in the sequel, we will define an abbreviation for the previous operation. We will denote ultimately :

$$\exp \left(- \langle \cdot \rangle_{\alpha_1 \rightarrow \alpha_2}^S \right) \equiv T \exp \left(- \int_{\alpha_2}^{\alpha_1} S_{\mathbf{X}_{\alpha_1}} \langle \cdot \rangle_\alpha S_{\mathbf{X}_\alpha} d\alpha \right) \quad (24)$$

as the exponentiated bracket plus shift, called *EBS* operation.

We can now exploit the mathematical construction developed in this section to consider the formal diagonalization of an arbitrary matrix valued quantum Hamiltonian.

III. THE DIAGONALIZATION PROCEDURE

We now consider a generic matrix valued quantum Hamiltonian $H(\mathbf{R}, \mathbf{P})$ where \mathbf{R} and \mathbf{P} are the usual canonical coordinate and momentum operators satisfying the canonical Heisenberg algebra. Our goal is to find a unitary transformation U such that UHU^+ is a diagonal matrix valued operator $\varepsilon(\mathbf{R}, \mathbf{P})$ (block diagonal for the Dirac Hamiltonian). This is in general an excessively difficult mathematical problem. For this reason we consider this problem by dividing it in several steps.

A. The Hamiltonian

First we introduce the unitary matrix $U_\alpha(\mathbf{X}_\alpha) \equiv U(\mathbf{X}_\alpha, \alpha)$ which diagonalizes the Hamiltonian $H(\mathbf{X}_\alpha)$ where the canonical variables $\mathbf{X}_\hbar \equiv \mathbf{X}$ have been replaced by the running ones \mathbf{X}_α , so that we can write

$$U_\alpha(\mathbf{X}_\alpha) H(\mathbf{X}_\alpha) U_\alpha^+(\mathbf{X}_\alpha) = \varepsilon_\alpha(\mathbf{X}_\alpha) \quad (25)$$

Note that following [5] we choose to put $U_\alpha(\mathbf{X}_\alpha)$ and $H(\mathbf{X}_\alpha)$ in a symmetric form. With the help of the previous identities Eqs. (9) (13) (14), we can compute $\varepsilon_\hbar(\mathbf{X}_\hbar) \equiv \varepsilon(\mathbf{X})$. Indeed as $\varepsilon(\mathbf{X}) = \mathcal{E}(\varepsilon(\mathbf{X}))$ we can write

$$\varepsilon(\mathbf{X}) = \mathcal{E} \left(\varepsilon_0(\mathbf{X}_0) + \int_0^{\hbar} d\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) \right). \quad (26)$$

Clearly $\varepsilon_0(\mathbf{X}_0)$ corresponds to the diagonal representation of the original Hamiltonian $H(\mathbf{X})$ where the canonical operators \mathbf{X} have been replaced by the classical variables \mathbf{X}_0 . In practice, it is usually quite easy to diagonalize the Hamiltonian when the operators are commuting. This is the essence of the method. Starting with classical variables we can recursively introduce more and more "quantification" through the running parameter α until we get the full quantum Hamiltonian. This procedure is now described in the following.

The quantity $\mathcal{E}(d\varepsilon_{\alpha_1}) = \mathcal{E}\left(\left(\frac{\partial}{\partial\alpha_1} + \langle \cdot \rangle\right)\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1})\right)$ can be straightforwardly computed by using the rule previously given for the bracket of a product Eq. (12). Indeed we find

$$\begin{aligned} \left(\frac{\partial}{\partial\alpha_1} + \langle \cdot \rangle\right)\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) &= \partial_{\alpha_1}U_{\alpha_1}U_{\alpha_1}^+\varepsilon_{\alpha_1} + \varepsilon_{\alpha_1}U_{\alpha_1}\partial_{\alpha_1}U_{\alpha_1}^+ + U_{\alpha_1}\frac{\partial}{\partial\alpha_1}H(\mathbf{X}_{\alpha_1})U_{\alpha_1}^+ \\ &\quad + \langle U_{\alpha_1} \rangle U_{\alpha_1}^+\varepsilon_{\alpha_1} + U_{\alpha_1}\langle H(\mathbf{X}_{\alpha_1}) \rangle U_{\alpha_1}^+ + \varepsilon_{\alpha_1}U_{\alpha_1}\langle U_{\alpha_1}^+ \rangle \\ &\quad - \frac{i}{2}(\nabla_{P_i}U_{\alpha_1}\nabla_{R_i}H(\mathbf{X}_{\alpha_1})U_{\alpha_1}^+ - \nabla_{R_i}U_{\alpha_1}\nabla_{P_i}H(\mathbf{X}_{\alpha_1})U_{\alpha_1}^+) \\ &\quad - \frac{i}{2}(U_{\alpha_1}\nabla_{P_i}H(\mathbf{X}_{\alpha_1})\nabla_{R_i}U_{\alpha_1}^+ - U_{\alpha_1}\nabla_{R_i}H(\mathbf{X}_{\alpha_1})\nabla_{P_i}U_{\alpha_1}^+) \\ &\quad - \frac{i}{2}(\nabla_{P_i}U_{\alpha_1}H(\mathbf{X}_{\alpha_1})\nabla_{R_i}U_{\alpha_1}^+ - \nabla_{R_i}U_{\alpha_1}H(\mathbf{X}_{\alpha_1})\nabla_{P_i}U_{\alpha_1}^+) \end{aligned} \quad (27)$$

We note that by construction $d\varepsilon_{\alpha_1}$ is a diagonal matrix so that we obviously have the following identities

$$\left(\frac{\partial}{\partial\alpha_1} + \langle \cdot \rangle\right)\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) = \mathcal{P}_+(\text{R.H.S. of Eq. 27}) \quad (28)$$

$$0 = \mathcal{P}_-(\text{R.H.S. of Eq. 27}) \quad (29)$$

where \mathcal{P}_+ and \mathcal{P}_- are the projection on the diagonal and off the diagonal respectively. As in [5] we now introduce the notations $\mathcal{A}_{\alpha_1}^{R_i} = iU_{\alpha_1}(\mathbf{X}_{\alpha_1})\nabla_{P_i}U_{\alpha_1}^+(\mathbf{X}_{\alpha_1})$ and $\mathcal{A}_{\alpha_1}^{P_i} = -iU_{\alpha_1}(\mathbf{X}_{\alpha_1})\nabla_{R_i}U_{\alpha_1}^+(\mathbf{X}_{\alpha_1})$ as well as $\mathcal{B}_{\alpha_1} = (\langle U_{\alpha_1}(\mathbf{X}_{\alpha_1}) \rangle)U_{\alpha_1}^+(\mathbf{X}_{\alpha_1})$ so that Eq. (28) can be written

$$\begin{aligned} \left(\frac{\partial}{\partial\alpha_1} + \langle \cdot \rangle\right)\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) &= \mathcal{P}_+\left(\partial_{\alpha_1}U_{\alpha_1}U_{\alpha_1}^+\varepsilon_{\alpha_1} + \varepsilon_{\alpha_1}U_{\alpha_1}\partial_{\alpha_1}U_{\alpha_1}^+ + U_{\alpha_1}\frac{\partial}{\partial\alpha_1}H(\mathbf{X}_{\alpha_1})U_{\alpha_1}^+\right) \\ &\quad + \mathcal{P}_+\left\{\frac{1}{2}\mathcal{A}_{\alpha_1}^{R_i}\nabla_{R_i}\varepsilon_{\alpha_1} + \nabla_{R_i}\varepsilon_{\alpha_1}\mathcal{A}_{\alpha_1}^{R_i} + \mathcal{A}_{\alpha_1}^{P_i}\nabla_{P_i}\varepsilon_{\alpha_1} + \nabla_{P_i}\varepsilon_{\alpha_1}\mathcal{A}_{\alpha_1}^{P_i}\right\} \\ &\quad + \frac{i}{2}\mathcal{P}_+\left\{\mathcal{A}_{\alpha_1}^{P_i}\varepsilon_{\alpha_1}\mathcal{A}_{\alpha_1}^{R_i} - \mathcal{A}_{\alpha_1}^{R_i}\varepsilon_{\alpha_1}\mathcal{A}_{\alpha_1}^{P_i} + \varepsilon_{\alpha_1}[\mathcal{A}_{\alpha_1}^{R_i}, \mathcal{A}_{\alpha_1}^{P_i}] + [\mathcal{A}_{\alpha_1}^{R_i}, \mathcal{A}_{\alpha_1}^{P_i}]\varepsilon_{\alpha_1}\right\} \\ &\quad + \mathcal{P}_+\left\{U_{\alpha_1}\langle H(\mathbf{X}_{\alpha_1}) \rangle U_{\alpha_1}^+ + [\mathcal{B}_{\alpha_1}\varepsilon_{\alpha_1} - \varepsilon_{\alpha_1}\mathcal{B}_{\alpha_1}^+]\right\}. \end{aligned} \quad (30)$$

This action is of affine type in $\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1})$ and will be denoted

$$\left(\frac{\partial}{\partial\alpha_1} + \langle \cdot \rangle\right)\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) = O_{\alpha_1}\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) \quad (31)$$

with O_{α_1} given by the r.h.s. expression of Eq. (30). Let us remark that this operator is well defined since, as explained before, the operation $\left(\frac{\partial}{\partial\alpha_1} + \langle \cdot \rangle\right)$ is independent from any symmetrization scheme for $\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1})$.

Now similarly to Eq. (26) we can write :

$$\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) = \varepsilon_0(\mathbf{X}_0) + \int_0^{\alpha_1} d\varepsilon_{\alpha_2}(\mathbf{X}_{\alpha_2}) = \varepsilon_0(\mathbf{X}_0) + \int_0^{\alpha_1} d(U(\mathbf{X}_{\alpha_2})H(\mathbf{X}_{\alpha_2})U^+(\mathbf{X}_{\alpha_2})) \quad (32)$$

which can be inserted in the expectation Eq. (26) to get the full quantum diagonal representation $\varepsilon(\mathbf{X}) \equiv \varepsilon_{\hbar}(\mathbf{X}_{\hbar})$ as :

$$\begin{aligned} \varepsilon_{\hbar}(\mathbf{X}_{\hbar}) &= \mathcal{E} \left(\varepsilon_0(\mathbf{X}_0) + \int_0^{\hbar} d\varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) \right) \\ &= \mathcal{E} \left(\varepsilon_0(\mathbf{X}_0) + \int_0^{\hbar} O_{\alpha_1} \cdot \left[\varepsilon_0(\mathbf{X}_0) + \int_0^{\alpha_1} d[U(\mathbf{X}_{\alpha_2})H(\mathbf{X}_{\alpha_2})U^+(\mathbf{X}_{\alpha_2})] \right] d\alpha_1 \right) \end{aligned} \quad (33)$$

Now remark that since $\alpha_2 < \alpha_1$, the terms $\nabla_{\mathbf{X}_{\alpha_1}} \varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) d\mathbf{X}_{\alpha_1}$ and $\nabla_{X_{\alpha_1}^j} \nabla_{X_{\alpha_1}^i} \varepsilon_{\alpha_1}(\mathbf{X}_{\alpha_1}) (dX_{\alpha_1}^i dX_{\alpha_1}^j + dX_{\alpha_1}^j dX_{\alpha_1}^i)$ in $d\varepsilon(\mathbf{X}_{\alpha_1})$ do not recombine with anything to give a product of the type $d\mathbf{R}_{\alpha_1} d\mathbf{P}_{\alpha_1}$ that would induce a $d\alpha_1$ contribution. They will thus cancel in the expectation. As a consequence, Eq. (33) can be written:

$$\varepsilon_{\hbar}(\mathbf{X}_{\hbar}) = \mathcal{E} \left(\varepsilon_0(\mathbf{X}_0) + \int_0^{\hbar} O_{\alpha_1} \cdot \left[\varepsilon_0(\mathbf{X}_0) + \int_0^{\alpha_1} O_{\alpha_2} \cdot \varepsilon_{\alpha_2}(\mathbf{X}_{\alpha_2}) d\alpha_2 \right] d\alpha_1 \right)$$

Repeating the procedure one can then show by iteration that :

$$\varepsilon_{\hbar}(\mathbf{X}_{\hbar}) = \mathcal{E} \left(\int_0^{\hbar} \left[\sum_{n=0}^{\infty} \int_{0 < \alpha_n < \dots < \alpha_1 < \alpha} O_{\alpha_1} \dots O_{\alpha_n} d\alpha_1 \dots d\alpha_n \right] \varepsilon_0(\mathbf{X}_0) \right) \quad (34)$$

At this point, an important comment has to be made. The operator O_{α_i} have been designed to depend on \mathbf{X}_{α_i} . We can use the connection operation defined in Eq. (24) to put the series appearing in the previous formula Eq. (34) at the same point. Once again, since $0 < \alpha_n < \dots < \alpha_1$, we can write, inside the expectation :

$$\begin{aligned} &\left[\int_{0 < \alpha_n < \dots < \alpha_1 < \hbar} O_{\alpha_1} \dots O_{\alpha_n} d\alpha_1 \dots d\alpha_n \right] \varepsilon_0(\mathbf{X}_0) \\ &= \left[\int_{0 < \alpha_n < \dots < \alpha_1 < \hbar} e^{-\langle \cdot \rangle_{\alpha_1 \rightarrow \hbar}^S} O_{\alpha_1} \dots e^{-\langle \cdot \rangle_{\alpha_{n-1} \rightarrow \alpha_{n-2}}^S} O_{\alpha_{n-1}} e^{-\langle \cdot \rangle_{\alpha_n \rightarrow \alpha_{n-1}}^S} O_{\alpha_n} e^{-\langle \cdot \rangle_{0 \rightarrow \alpha_n}^S} \right] \varepsilon_0(\mathbf{X}_0) \end{aligned}$$

The computation of this last quantity derives directly from the definitions of the various operators involved and proceeds as follows :

starting from $\varepsilon_0(\mathbf{X}_0)$, first apply $e^{-\langle \cdot \rangle_{0 \rightarrow \alpha_n}^S}$ on $\varepsilon_0(\mathbf{X}_0)$ that will shift \mathbf{X}_0 by \mathbf{X}_{α_n} . Then make O_{α_n} acts on the result, apply $e^{-\langle \cdot \rangle_{\alpha_n \rightarrow \alpha_{n-1}}^S}$ which replaces \mathbf{X}_{α_n} by $\mathbf{X}_{\alpha_{n-1}}$ and so on. Ultimately at the end of the process, \mathbf{X}_0 is replaced by \mathbf{X}_{\hbar} , this last variable being independent from the integration variables, the various integrals can be computed easily, as the

integral of a series expansion. This process of integration will be of course the same for all similar expressions in the sequel.

We can now deduce that (with the notation $\mathbf{X} \equiv \mathbf{X}_{\hbar}$):

$$\varepsilon(\mathbf{X}) \equiv \varepsilon_{\hbar}(\mathbf{X}_{\hbar}) = \mathcal{E} \left(\left[\mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} e^{-\langle \cdot \rangle_{\alpha \rightarrow \hbar}^S} \exp(\alpha \langle \cdot \rangle) O_{\alpha} e^{-\langle \cdot \rangle_{0 \rightarrow \alpha}^S} d\alpha \right] \right] \varepsilon_0(\mathbf{X}_0) \right) \quad (35)$$

which is compact expression for the diagonal Hamiltonian $\varepsilon(\mathbf{R}, \mathbf{P})$ in terms of the "classical" diagonal Hamiltonian $\varepsilon_0(\mathbf{R}, \mathbf{P})$ in which the classical variables $\mathbf{R}_0, \mathbf{P}_0$ have been now replaced by the quantum ones \mathbf{R}, \mathbf{P} due to the *EBS* action. Here \mathcal{T} is the usual notation for the "time ordered product". Eq. (35) is the required expression and constitutes the main result of this paper.

Clearly the practical application of Eq. (35) requires the knowledge of the transformation matrices U_{α} which enter into the definition of the operators O_{α} .

B. The transformation matrix U

Note first that there is a certain arbitrariness in the choice of the unitary matrix $U_{\alpha}(\mathbf{X})$ as explained in [5] which reflects a kind of gauge invariance. Actually, multiplying the transformation matrix $U_{\alpha}(\mathbf{X})$ on the right by a diagonal unitary matrix yields an other diagonalization, equivalent to the previous one. In particular it allows to choose n conditions for the diagonal entries of $U_{\alpha}(\mathbf{X})$ ($n \times n$ being the size of $U_{\alpha}(\mathbf{X})$). An explicit choice will be done below to simplify our expressions.

To find the transformation matrix U we use the same approach as for the diagonalization of the Hamiltonian, by writing :

$$U_{\hbar}(\mathbf{X}_{\hbar}) = \mathcal{E} \left(U(\mathbf{X}_0) + \int_0^{\hbar} dU_{\alpha}(\mathbf{X}_{\alpha}) \right) \quad (36)$$

$$= \mathcal{E} \left(\int_0^{\hbar} \left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) U_{\alpha}(\mathbf{X}_{\alpha}) d\alpha \right) \quad (37)$$

we can find $(\partial_{\alpha} + \langle \cdot \rangle) U_{\alpha}(\mathbf{X}_{\alpha})$ by using again the diagonalization process leading to Eq.

(27). Indeed we have

$$\begin{aligned}
\left(\frac{\partial}{\partial\alpha} + \langle.\rangle\right) \varepsilon_\alpha(\mathbf{X}_\alpha) &= \left(\left(\frac{\partial}{\partial\alpha} + \langle.\rangle\right) U_\alpha\right) U_\alpha^+ \varepsilon_\alpha + \varepsilon_\alpha U_\alpha \left(\frac{\partial}{\partial\alpha} + \langle.\rangle\right) U_\alpha^+ \\
&+ U_\alpha \left(\left(\frac{\partial}{\partial\alpha} + \langle.\rangle\right) H(\mathbf{X}_\alpha)\right) U_\alpha^+ \\
&- \frac{i}{2} (\nabla_{P_i} U_\alpha \nabla_{R_i} H(\mathbf{X}_\alpha) U_\alpha^+ - \nabla_{R_i} U_\alpha \nabla_{P_i} H(\mathbf{X}_\alpha) U_\alpha^+) \\
&- \frac{i}{2} (U_\alpha \nabla_{P_i} H(\mathbf{X}_\alpha) \nabla_{R_i} U_\alpha^+ - U_\alpha \nabla_{R_i} H(\mathbf{X}_\alpha) \nabla_{P_i} U_\alpha^+) \\
&- \frac{i}{2} (\nabla_{P_i} U_\alpha H(\mathbf{X}_\alpha) \nabla_{R_i} U_\alpha^+ - \nabla_{R_i} U_\alpha H(\mathbf{X}_\alpha) \nabla_{P_i} U_\alpha^+) \quad (38)
\end{aligned}$$

In addition we have from the unitarity condition $U_\alpha U_\alpha^+ = 1$ the relation $(\partial_\alpha + \langle.\rangle)(U_\alpha U_\alpha^+) = 0$ which reads :

$$0 = \left(\left(\frac{\partial}{\partial\alpha} + \langle.\rangle\right) U_\alpha\right) U_\alpha^+ + U_\alpha \left(\frac{\partial}{\partial\alpha} + \langle.\rangle\right) U_\alpha^+ - \frac{i}{2} (\nabla_{P_i} U_\alpha \nabla_{R_i} U_\alpha^+ - \nabla_{R_i} U_\alpha \nabla_{P_i} U_\alpha^+) \quad (39)$$

Mixing the two equations Eqs. (38)(39), we obtain that Eq. (38) after projection on the non diagonal part becomes the equality

$$\begin{aligned}
\mathcal{P}_- \left[\left(\left(\frac{\partial}{\partial\alpha} + \langle.\rangle\right) U_\alpha\right) U_\alpha^+, \varepsilon_\alpha\right] &= -\mathcal{P}_- \left(U_\alpha \left(\left(\frac{\partial}{\partial\alpha} + \langle.\rangle\right) H(\mathbf{X}_\alpha)\right) U_\alpha^+\right) \\
&- \frac{i}{2} \varepsilon_\alpha \mathcal{P}_- (\nabla_{P_i} U_\alpha \nabla_{R_i} U_\alpha^+ - \nabla_{R_i} U_\alpha \nabla_{P_i} U_\alpha^+) \\
&+ \frac{i}{2} \mathcal{P}_- (\nabla_{P_i} U_\alpha \nabla_{R_i} H(\mathbf{X}_\alpha) U_\alpha^+ - \nabla_{R_i} U_\alpha \nabla_{P_i} H(\mathbf{X}_\alpha) U_\alpha^+) \\
&+ \frac{i}{2} \mathcal{P}_- (U_\alpha \nabla_{P_i} H(\mathbf{X}_\alpha) \nabla_{R_i} U_\alpha^+ - U_\alpha \nabla_{R_i} H(\mathbf{X}_\alpha) \nabla_{P_i} U_\alpha^+) \\
&+ \frac{i}{2} \mathcal{P}_- (\nabla_{P_i} U_\alpha H(\mathbf{X}_\alpha) \nabla_{R_i} U_\alpha^+ - \nabla_{R_i} U_\alpha H(\mathbf{X}_\alpha) \nabla_{P_i} U_\alpha^+) \quad (40)
\end{aligned}$$

Both conditions Eqs. (39)(40) can be solved for $(\partial_\alpha + \langle.\rangle)U_\alpha$. Actually, decomposing $[(\partial_\alpha + \langle.\rangle)U_\alpha]U_\alpha^+ = hr + ahr$ in Hermitian hr and anti-Hermitian ahr part, Eq. (40) reads

$$\mathcal{P}_- [hr + ahr, \varepsilon_\alpha] = M(U_\alpha) \quad (41)$$

where $M(U_\alpha)$ is the r.h.s of Eq. (40). We can now use a particular choice of gauge. We fix our diagonalization process by setting $P_+(ahr) = 0$. Practically it corresponds to multiply U_α on the right by a unitary diagonal matrix D_α (which has n degree of freedom) such that $P_+\left(\left(\frac{\partial}{\partial\alpha} + \langle.\rangle\right)U_\alpha D_\alpha\right)D_\alpha^+ U_\alpha^+$ is Hermitian. This gives n conditions that determine D_α .

Now, from Eq. (41) and given our gauge choice for U_α , we deduce that $\mathcal{P}_-(ahr) = ahr$ and $[ahr, \varepsilon_\alpha] = \frac{1}{2}(M(U_\alpha) + M^+(U_\alpha))$. Then we can formally write ahr under the form

$$ahr = \tilde{N}(\alpha) \cdot U_\alpha \quad (42)$$

where we have defined :

$$\tilde{N}(\alpha) \cdot U_\alpha = \frac{1}{2}([\cdot, \varepsilon_\alpha]^{-1})(M(U_\alpha) + M^+(U_\alpha)) \quad (43)$$

and the inverse of the commutator operation $[\cdot, \varepsilon_\alpha]$ has the following meaning

$$\begin{aligned} [[\cdot, \varepsilon_\alpha]^{-1} \cdot M, \varepsilon_\alpha] &= [\cdot, \varepsilon_\alpha]^{-1} \cdot [M, \varepsilon_\alpha] = M \quad \text{for} \quad [M, \varepsilon_\alpha] \neq 0 \\ [\cdot, \varepsilon_\alpha]^{-1} \cdot M &= 0 \quad \text{if} \quad [M, \varepsilon_\alpha] = 0 \end{aligned}$$

To complete the determination of U_α we still need to deduce hr which is readily obtained from the unitarity condition Eq. (39) as:

$$hr = -\frac{i}{4}(\nabla_{P_i} U_\alpha \nabla_{R_i} U_\alpha^+ - \nabla_{R_i} U_\alpha \nabla_{P_i} U_\alpha^+) \quad (44)$$

Gathering the results Eqs.(43) (44) allows us to introduce the operator N_α operating on U_α in the following manner $N_\alpha \cdot U_\alpha = (hr + ahr) U_\alpha$, so that we can write

$$\left[\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) U_\alpha \right] = N_\alpha \cdot U_\alpha$$

with N_α given explicitly by the following expression

$$N_\alpha \cdot U_\alpha = \frac{1}{2}([\cdot, \varepsilon_\alpha]^{-1})(M(U_\alpha) + M^+(U_\alpha)) U_\alpha + \frac{i}{4}(\nabla_{P_i} U_\alpha \nabla_{R_i} U_\alpha^+ - \nabla_{R_i} U_\alpha \nabla_{P_i} U_\alpha^+) U_\alpha$$

As for the energy diagonalization, this expression can be rewritten in terms of physical quantities

$$\begin{aligned} N_\alpha \cdot U_\alpha &= -[\cdot, \varepsilon_\alpha]^{-1} \cdot \left[\mathcal{P}_- \left\{ \frac{1}{2} \mathcal{A}_\alpha^{R_i} \nabla_{R_i} \varepsilon_\alpha + \nabla_{R_i} \varepsilon_\alpha \mathcal{A}_\alpha^{R_i} + \mathcal{A}_\alpha^{P_i} \nabla_{P_i} \varepsilon_\alpha + \nabla_{P_i} \varepsilon_\alpha \mathcal{A}_\alpha^{P_i} \right\} \right. \\ &\quad \left. + \frac{i}{4} \mathcal{P}_- \{ [\varepsilon_\alpha, \mathcal{A}_\alpha^{R_i}] \mathcal{A}_\alpha^{P_i} - [\varepsilon_\alpha, \mathcal{A}_\alpha^{P_i}] \mathcal{A}_\alpha^{R_i} \} + H.C. \right] - \frac{i}{4} \mathcal{P}_- \{ [\mathcal{A}_\alpha^{R_i}, \mathcal{A}_\alpha^{P_i}] U_\alpha \} \quad (45) \end{aligned}$$

where $H.C.$ stands for the Hermitic conjugate of the right hand side and $\varepsilon_\alpha(\mathbf{X}_\alpha)$ is computed recursively as explained before.

As for $\varepsilon(\mathbf{X})$, we can therefore write for $U(\mathbf{X}) \equiv U_{\hbar}(\mathbf{X}_{\hbar})$:

$$U(\mathbf{X}) = \mathcal{E} \left(\left[\mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} e^{-\langle \cdot \rangle_{\alpha \rightarrow \hbar}^S} N_{\alpha} \exp e^{-\langle \cdot \rangle_{0 \rightarrow \alpha}^S} d\alpha \right] \right] U_0(\mathbf{X}_0) \right) \quad (46)$$

This expression has a very similar structure as the solution for the energy Eq. (35) except that the operator O_{α} has to be replaced by N_{α} .

Note that given the solution derived for U_{α} we can now rewrite the operator O_{α} in an simpler form. Indeed starting again from

$$\begin{aligned} O_{\alpha} &= \left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) \varepsilon_{\alpha}(\mathbf{X}_{\alpha}) = \mathcal{P}_+ \left[\left(\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) U_{\alpha} \right) U_{\alpha}^+ \varepsilon_{\alpha} + \varepsilon_{\alpha} U_{\alpha} \left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) U_{\alpha}^+ \right] \\ &+ \mathcal{P}_+ \left[U_{\alpha} \left(\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) H(\mathbf{X}_{\alpha}) \right) U_{\alpha}^+ \right] - \frac{i}{2} \mathcal{P}_+ (\nabla_{P_i} U_{\alpha} \nabla_{R_i} H(\mathbf{X}_{\alpha}) U_{\alpha}^+ - \nabla_{R_i} U_{\alpha} \nabla_{P_i} H(\mathbf{X}_{\alpha}) U_{\alpha}^+) \\ &- \frac{i}{2} \mathcal{P}_+ (U_{\alpha} \nabla_{P_i} H(\mathbf{X}_{\alpha}) \nabla_{R_i} U_{\alpha}^+ - U_{\alpha} \nabla_{R_i} H(\mathbf{X}_{\alpha}) \nabla_{P_i} U_{\alpha}^+) \\ &- \frac{i}{2} \mathcal{P}_+ (\nabla_{P_i} U_{\alpha} H(\mathbf{X}_{\alpha}) \nabla_{R_i} U_{\alpha}^+ - \nabla_{R_i} U_{\alpha} H(\mathbf{X}_{\alpha}) \nabla_{P_i} U_{\alpha}^+) \end{aligned} \quad (47)$$

using the unitarity condition for U_{α} yields :

$$\begin{aligned} O_{\alpha} \varepsilon_{\alpha}(\mathbf{X}_{\alpha}) &= \mathcal{P}_+ \left[\left[\left(\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) U_{\alpha} \right) U_{\alpha}^+, \varepsilon_{\alpha} \right] + U_{\alpha} \left(\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) H(\mathbf{X}_{\alpha}) \right) U_{\alpha}^+ \right] \\ &+ \frac{i}{2} \varepsilon_{\alpha} \mathcal{P}_+ (\nabla_{P_i} U_{\alpha} \nabla_{R_i} U_{\alpha}^+ - \nabla_{R_i} U_{\alpha} \nabla_{P_i} U_{\alpha}^+) \\ &- \frac{i}{2} \mathcal{P}_+ (\nabla_{P_i} U_{\alpha} \nabla_{R_i} H(\mathbf{X}_{\alpha}) U_{\alpha}^+ - \nabla_{R_i} U_{\alpha} \nabla_{P_i} H(\mathbf{X}_{\alpha}) U_{\alpha}^+) \\ &- \frac{i}{2} \mathcal{P}_+ (U_{\alpha} \nabla_{P_i} H(\mathbf{X}_{\alpha}) \nabla_{R_i} U_{\alpha}^+ - U_{\alpha} \nabla_{R_i} H(\mathbf{X}_{\alpha}) \nabla_{P_i} U_{\alpha}^+) \\ &- \frac{i}{2} \mathcal{P}_+ (\nabla_{P_i} U_{\alpha} H(\mathbf{X}_{\alpha}) \nabla_{R_i} U_{\alpha}^+ - \nabla_{R_i} U_{\alpha} H(\mathbf{X}_{\alpha}) \nabla_{P_i} U_{\alpha}^+) \end{aligned} \quad (48)$$

Now, using our gauge condition which states that the antihermitian part of $\mathcal{P}_+ \left(\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) U_{\alpha} \right) U_{\alpha}^+$ is null, as well as the fact that $\varepsilon_{\alpha}(\mathbf{X}_{\alpha})$ is assumed to be an Hermitian operator, we are led to the expression :

$$\begin{aligned} 2O_{\alpha} \varepsilon_{\alpha}(\mathbf{X}_{\alpha}) &= \mathcal{P}_+ \left[U_{\alpha} \left(\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) H(\mathbf{X}_{\alpha}) \right) U_{\alpha}^+ \right] \\ &+ \frac{i}{2} \varepsilon_{\alpha} \mathcal{P}_+ (\nabla_{P_i} U_{\alpha} \nabla_{R_i} U_{\alpha}^+ - \nabla_{R_i} U_{\alpha} \nabla_{P_i} U_{\alpha}^+) \\ &- \frac{i}{2} \mathcal{P}_+ (\nabla_{P_i} U_{\alpha} \nabla_{R_i} H(\mathbf{X}_{\alpha}) U_{\alpha}^+ - \nabla_{R_i} U_{\alpha} \nabla_{P_i} H(\mathbf{X}_{\alpha}) U_{\alpha}^+) \\ &- \frac{i}{2} \mathcal{P}_+ (U_{\alpha} \nabla_{P_i} H(\mathbf{X}_{\alpha}) \nabla_{R_i} U_{\alpha}^+ - U_{\alpha} \nabla_{R_i} H(\mathbf{X}_{\alpha}) \nabla_{P_i} U_{\alpha}^+) \\ &- \frac{i}{2} \mathcal{P}_+ (\nabla_{P_i} U_{\alpha} H(\mathbf{X}_{\alpha}) \nabla_{R_i} U_{\alpha}^+ - \nabla_{R_i} U_{\alpha} H(\mathbf{X}_{\alpha}) \nabla_{P_i} U_{\alpha}^+) + H.C. \end{aligned} \quad (49)$$

This can be rewritten in terms of physical operators :

$$\begin{aligned}
O_\alpha \varepsilon_\alpha(\mathbf{X}_\alpha) &= \frac{1}{2} \mathcal{P}_+ \left\{ \mathcal{A}_\alpha^{R_l} \nabla_{R_l} \varepsilon_\alpha + \nabla_{R_l} \varepsilon_\alpha \mathcal{A}_\alpha^{R_l} + \mathcal{A}_\alpha^{P_l} \nabla_{P_l} \varepsilon_\alpha + \nabla_{P_l} \varepsilon_\alpha \mathcal{A}_\alpha^{P_l} \right\} \\
&+ \left[\frac{i}{4} \mathcal{P}_+ \left\{ [\varepsilon_\alpha, \mathcal{A}_\alpha^{R_l}] \mathcal{A}_\alpha^{P_l} - [\varepsilon_\alpha, \mathcal{A}_\alpha^{P_l}] \mathcal{A}_\alpha^{R_l} \right\} + H.C. \right] + \mathcal{P}_+ \left[U_\alpha \left(\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) H(\mathbf{X}_\alpha) \right) U_\alpha^+ \right]
\end{aligned} \tag{50}$$

Let us remark, that in most of the physical applications (in particular for Dirac and Bloch electrons) the term $\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) H(\mathbf{X}_\alpha)$ will cancel.

C. The full system \mathbf{H} and \mathbf{U}

We can now write the solution of our diagonalization procedure for a general matrix valued Hamiltonian through an unitary transformation U as the solution of the following system of differential equations

$$\varepsilon(\mathbf{X}) = \mathcal{E} \left(\left[\mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} e^{-\langle \cdot \rangle_\alpha^S / \hbar} O_\alpha e^{-\langle \cdot \rangle_{0 \rightarrow \alpha}^S} d\alpha \right] \varepsilon_0(\mathbf{X}_0) \right] \right) \tag{51}$$

$$U(\mathbf{X}) = \mathcal{E} \left(\left[\mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} e^{-\langle \cdot \rangle_\alpha^S / \hbar} N_\alpha e^{-\langle \cdot \rangle_{0 \rightarrow \alpha}^S} d\alpha \right] U_0(\mathbf{X}_0) \right] \right) \tag{52}$$

where O_α and N_α are given respectively by Eqs.(50) and (45). The only pre-requirement is that the diagonal form at $\alpha = 0$, $\varepsilon_0(\mathbf{R}_0, \mathbf{P}_0)$ is known, i.e., when \mathbf{R} and \mathbf{P} are considered as classical commuting variables (which means that $U_0(\mathbf{X}_0)$ is known). Of course these equations do not allow to find directly $\varepsilon(\mathbf{X})$, $U(\mathbf{X})$ since those quantities are involved on the R.H.S. of these relations. However, they allow us to produce solutions for $\varepsilon(\mathbf{X})$ and $U(\mathbf{X})$ recursively in a series expansion in \hbar . Moreover, and as needed, the results of our previous sections show that the matrices $\varepsilon(\mathbf{X})$ and $U(\mathbf{X})$ obtained through this process are independent of any choice of symmetrization. A particular choice in the way of arranging the variables for $U(\mathbf{X}) = 0$ will lead to differently symmetrized, but identical, operators.

As it will appear clearly later on, having both $\varepsilon(\mathbf{X})$ and $U(\mathbf{X})$ at order n in \hbar , and reinserting in the exponential of Eqs. (51) (52) allows us to find $\varepsilon(\mathbf{X})$ and $U(\mathbf{X})$ at order $n+1$ in \hbar . But before solving recursively the set of equations Eqs.(51) (52), we first compare the present approach with the one developed in article [5]

IV. LINK WITH THE DIFFERENTIAL EQUATION OF REF. [5]

In [5] we developed a different less general approach which led to the differential equation

$$\begin{aligned} \frac{d}{d\alpha}\varepsilon_\alpha(\mathbf{X}_\alpha) &= [\partial_\alpha U_\alpha(\mathbf{X}_\alpha) U_\alpha^+(\mathbf{X}_\alpha), \varepsilon_\alpha(\mathbf{X}_\alpha)] + \frac{1}{2} \{ \mathcal{A}_\alpha^{Rl} \nabla_{Rl} \varepsilon_\alpha + \nabla_{Rl} \varepsilon_\alpha \mathcal{A}_\alpha^{Rl} + \mathcal{A}_\alpha^{Pl} \nabla_{Pl} \varepsilon_\alpha + \nabla_{Pl} \varepsilon_\alpha \mathcal{A}_\alpha^{Pl} \} \\ &+ \frac{i}{2} \{ [\varepsilon_\alpha, \mathcal{A}_\alpha^{Rl}] \mathcal{A}_\alpha^{Pl} - [\varepsilon_\alpha, \mathcal{A}_\alpha^{Pl}] \mathcal{A}_\alpha^{Rl} - [\varepsilon_\alpha, [\mathcal{A}_\alpha^{Rl}, \mathcal{A}_\alpha^{Pl}]] \} \\ &+ \{ U_\alpha \langle H(\mathbf{X}_\alpha) \rangle U_\alpha^+ + [\mathcal{B}_\alpha \varepsilon_\alpha - \varepsilon_\alpha \mathcal{B}_\alpha^+] - \langle \varepsilon_\alpha \rangle \} \end{aligned} \quad (53)$$

which was coupled to the evolution of the transformation matrix $U_\alpha(\mathbf{X}_\alpha)$ as a function of α [5] :

$$0 = \partial_\alpha U_\alpha(\mathbf{X}_\alpha) U_\alpha^+(\mathbf{X}_\alpha) + U_\alpha(\mathbf{X}_\alpha) \partial_\alpha U_\alpha^+(\mathbf{X}_\alpha) - \frac{i}{2} (\mathcal{B}_\alpha - \mathcal{B}_\alpha^+) - \frac{i}{2} [\mathcal{A}_\alpha^{Rl}, \mathcal{A}_\alpha^{Pl}] \quad (54)$$

where $\mathcal{B}_\alpha = (Asym[\nabla_{Rl} \nabla_{Pl} U_\alpha(\mathbf{X}_\alpha)]) U_\alpha^+(\mathbf{X}_\alpha)$. With these two equations Eqs. (53) and (54) at hand, the diagonalization process can be performed. Actually, since all quantities are matrix valued and since $\varepsilon_\alpha(\mathbf{X}_\alpha)$ is by definition a diagonal matrix, we can separate the energy equation Eq. (53) in a diagonal and a off-diagonal part such that we are led to the following two equations

$$\frac{d}{d\alpha}\varepsilon_\alpha(\mathbf{X}_\alpha) = \mathcal{P}_+[\text{R.H.S. of Eq. 53}] \quad (55)$$

$$0 = \mathcal{P}_-[\text{R.H.S. of Eq. 53}] \quad (56)$$

In [5] it was claimed that those three Eqs. (54) (55) (56) allow us to determine recursively in powers of α the energy of the quantum system in question. Actually, the integration over α of Eq. (55) gives $\varepsilon_\alpha(\mathbf{X}_\alpha)$ at order n in α when knowing all quantities at order $n-1$. By the same token, Eqs. (56) and (54) (whose meaning is that $U_\alpha(\mathbf{X}_\alpha)$ is unitary at each order in α) involve $\partial_\alpha U_\alpha(\mathbf{X}_\alpha)$, and allow to recover $U_\alpha(\mathbf{X}_\alpha)$ at order n by integration over α . As a consequence, the diagonalization process is perfectly controlled order by order in the series expansion in α . In [5] we also provided two physical examples at the order \hbar^2 .

Now we want to show that our solution Eq. (35) satisfies the differential equation Eq. (53). Indeed, from Eq. (26) we can write

$$\begin{aligned}
\frac{\partial}{\partial \hbar} \varepsilon(\mathbf{X}_{\hbar}) &= \mathcal{E} \left(\left[\mathcal{T} \exp \left[\int_{0 < \alpha_1 < \hbar} e^{-\langle \cdot \rangle_{\alpha_1 \rightarrow \hbar}^S} O_{\alpha_1} e^{-\langle \cdot \rangle_{0 \rightarrow \alpha_1}^S} d\alpha_1 \right] \right] \varepsilon_0(\mathbf{X}_0) \right) \\
&= \langle \cdot \rangle \varepsilon_{\hbar}(\mathbf{X}_{\hbar}) \\
\mathcal{E} \left(\sum_n \left[\int_{0 < \alpha_n < \dots < \alpha_2 < \hbar} O_{\hbar} e^{-\langle \cdot \rangle_{\alpha_2 \rightarrow \hbar}^S} O_{\alpha_2} \dots e^{-\langle \cdot \rangle_{\alpha_{n-1} \rightarrow \alpha_{n-2}}^S} O_{\alpha_{n-1}} e^{-\langle \cdot \rangle_{\alpha_n \rightarrow \alpha_{n-1}}^S} O_{\alpha_n} e^{-\langle \cdot \rangle_{0 \rightarrow \alpha_n}^S} \right] \varepsilon_0(\mathbf{X}_0) \right) \\
&= -\langle \cdot \rangle \varepsilon_{\hbar}(\mathbf{X}_{\hbar}) + \\
O_{\hbar} \mathcal{E} \left(\sum_n \left[\int_{0 < \alpha_n < \dots < \alpha_2 < \hbar} e^{-\langle \cdot \rangle_{\alpha_2 \rightarrow \hbar}^S} O_{\alpha_2} \dots e^{-\langle \cdot \rangle_{\alpha_{n-1} \rightarrow \alpha_{n-2}}^S} O_{\alpha_{n-1}} e^{-\langle \cdot \rangle_{\alpha_n \rightarrow \alpha_{n-1}}^S} O_{\alpha_n} e^{-\langle \cdot \rangle_{0 \rightarrow \alpha_n}^S} \right] \varepsilon_0(\mathbf{X}_0) \right) \\
&= -\langle \cdot \rangle \varepsilon_{\hbar}(\mathbf{X}_{\hbar}) + O_{\hbar} \varepsilon_{\hbar}(\mathbf{X}_{\hbar}) \tag{57}
\end{aligned}$$

The first equality has been obtained by using that by construction, $\frac{\partial}{\partial \hbar} e^{-\langle \cdot \rangle_{\alpha_1 \rightarrow \hbar}^S} = -S_{\mathbf{X}_{\hbar}} \langle \cdot \rangle_{\hbar} S_{\mathbf{X}_{\hbar}}$. We have given previously the expression for $O_{\hbar} \varepsilon_{\hbar}(\mathbf{Y}_{\hbar})$. Using moreover that $\langle \cdot \rangle \varepsilon_{\hbar}(\mathbf{X}_{\hbar}) = -\frac{i}{2} \text{Asym} \{ \nabla_{P_l} \nabla_{R_l} \varepsilon_{\alpha}(\mathbf{X}_{\alpha}) \}$ and Eq. (54) one gets directly that $(\varepsilon(\mathbf{R}, \mathbf{P}) - \varepsilon_0(\mathbf{R}, \mathbf{P}))$ is a solution of the differential (53).

V. DYNAMICAL OPERATORS AND COMMUTATION ALGEBRA

In this section we will see that new non-commuting position and momentum operators which have contributions from Berry connections emerge during the diagonalization and are more suitably to correspond to physical operators (for the physical discussion of this point see [11][7][5][12]).

From Eq.(27) one sees that the operator $O_{\alpha} = (\partial_{\alpha} + \langle \cdot \rangle)$ acting on $\varepsilon_0(\mathbf{X}_{\alpha})$ can be decomposed as a sum of a "translation" operator T and a "magnetization" M operator (this terminology is explained in [5])

$$O_{\alpha} \varepsilon_0(\mathbf{X}_{\alpha}) = (T_{\alpha} + M_{\alpha}) \varepsilon_0(\mathbf{X}_{\alpha}). \tag{58}$$

where the "magnetization" operator acts as

$$M_{\alpha} \varepsilon_0(\mathbf{X}_{\alpha}) = \frac{i}{2} \mathcal{P}_+ \{ [\varepsilon_{\alpha}, \mathcal{A}_{\alpha}^{R_i}] \mathcal{A}_{\alpha}^{P_i} - [\varepsilon_{\alpha}, \mathcal{A}_{\alpha}^{P_i}] \mathcal{A}_{\alpha}^{R_i} \} + \mathcal{P}_+ \left[U_{\alpha} \left(\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) H(\mathbf{X}_{\alpha}) \right) U_{\alpha}^+ \right] \tag{59}$$

and

$$T_\alpha \varepsilon_0(\mathbf{X}_\alpha) = \frac{1}{2} \mathcal{P}_+ \left\{ \mathcal{A}_\alpha^{R_l} \nabla_{R_l} \varepsilon_0(\mathbf{X}_\alpha) + \nabla_{R_l} \varepsilon_0 \mathcal{A}_\alpha^{R_l}(\mathbf{X}_\alpha) + \mathcal{A}_\alpha^{P_l} \nabla_{P_l} \varepsilon_0(\mathbf{X}_\alpha) + \nabla_{P_l} \varepsilon_0 \mathcal{A}_\alpha^{P_l}(\mathbf{X}_\alpha) \right\} \quad (60)$$

To inspect the action of T_α , let us skip the magnetization contribution M_α for the sake of clarity, and consider the following relevant contribution for the computation of the diagonalized Hamiltonian :

$$\mathcal{E} \mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} e^{-\langle \cdot \rangle_{\alpha \rightarrow \hbar}^S} T_\alpha e^{-\langle \cdot \rangle_{0 \rightarrow \alpha}^S} d\alpha \right] \varepsilon_0(\mathbf{X}_0)$$

When developed in series (as explained in the previous sections), it is given by

$$\mathcal{E} \sum_{n=0}^n \left[\int_{0 < \alpha_n < \dots < \alpha_1 < \hbar} e^{-\langle \cdot \rangle_{\alpha_1 \rightarrow \hbar}^S} T_{\alpha_1} \dots e^{-\langle \cdot \rangle_{\alpha_{n-1} \rightarrow \alpha_{n-2}}^S} T_{\alpha_{n-1}} e^{-\langle \cdot \rangle_{\alpha_n \rightarrow \alpha_{n-1}}^S} T_{\alpha_n} e^{-\langle \cdot \rangle_{0 \rightarrow \alpha_n}^S} \right] \varepsilon_0(\mathbf{X}_0) \quad (61)$$

Recall that each operation of the kind $e^{-\langle \cdot \rangle_{\alpha_i \rightarrow \alpha_{i-1}}^S}$ is going along with a change of variable $\mathbf{X}_{\alpha_i} \rightarrow \mathbf{X}_{\alpha_{i-1}}$.

We find $\mathcal{E} \mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} e^{-\langle \cdot \rangle_{\alpha \rightarrow \hbar}^S} T_\alpha e^{-\langle \cdot \rangle_{0 \rightarrow \alpha}^S} d\alpha \right] \varepsilon_0(\mathbf{X})$ by the following ansatz. We assume that $\mathcal{E} \mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} e^{-\langle \cdot \rangle_{\alpha \rightarrow \hbar}^S} T_\alpha e^{-\langle \cdot \rangle_{0 \rightarrow \alpha}^S} d\alpha \right] \varepsilon_0(\mathbf{X}) = \varepsilon_0(\mathbf{X}_\hbar + A_\hbar^{\mathbf{X}})$ where $A_\hbar^{\mathbf{X}} \equiv A_\hbar^{\mathbf{X}}(\mathbf{X}_\hbar, \hbar)$ has to be determined and a particular (arbitrary) choice of symmetrization of the variables \mathbf{X} has been made. Due to the form of the left hand side in the previous relation Eq. (61), $A_\hbar^{\mathbf{X}}$ is of order \hbar . We choose to write ε_0 such that the powers of \mathbf{R} and \mathbf{P} are in a completely symmetrized form, that is we sum over all equally weighted permutations of the canonical variables in the series expansion of $\varepsilon_0(\mathbf{X})$. Then, $\varepsilon_0(\mathbf{X} + A_\hbar^{\mathbf{X}})$ is obtained by replacing \mathbf{X} by $\mathbf{X} + A_\hbar^{\mathbf{X}}$ in the series expansion of ε_0 . Note that, for later purposes $\langle \cdot \rangle_\hbar \varepsilon_0(\mathbf{X} + A_\hbar^{\mathbf{X}})$ is of order \hbar^2 . To differentiate the right hand side of Eq. (61) with respect to \hbar we proceed in the following way. Decrease \hbar by an amount of $d\hbar$ and introduce the connection operation defined earlier so that $\mathcal{E} \varepsilon_0(\mathbf{X}_{\hbar-d\hbar} + A_{\hbar-d\hbar}^{\mathbf{X}}) = \mathcal{E} e^{-\langle \cdot \rangle_{\hbar-d\hbar \rightarrow \hbar}^S} \varepsilon_0(\mathbf{X}_{\hbar-d\hbar} + A_{\hbar-d\hbar}^{\mathbf{X}})$. As a consequence, the infinitesimal variation $\varepsilon_0(\mathbf{X}_\hbar + A_\hbar^{\mathbf{X}}) - \mathcal{E} \varepsilon_0(\mathbf{X}_{\hbar-d\hbar} + A_{\hbar-d\hbar}^{\mathbf{X}})$ is given by $\left[\frac{\partial A_\hbar^{\mathbf{X}}}{\partial \hbar} \nabla_{\mathbf{X}} \varepsilon_0(\mathbf{X} + A_\hbar^{\mathbf{X}}) + \langle \cdot \rangle_\hbar \varepsilon_0(\mathbf{X} + A_\hbar^{\mathbf{X}}) \right] d\hbar$. Note that in the last expression, the product $\frac{\partial A_\hbar^{\mathbf{X}}}{\partial \hbar} \nabla_{\mathbf{X}} \varepsilon_0(\mathbf{X} + A_\hbar^{\mathbf{X}})$ has to be understood in the sense that $\frac{dA_\hbar^{\mathbf{X}}}{d\hbar}$ replaces $\mathbf{X} + A_\hbar^{\mathbf{X}}$ at each place the gradient is acting on the series expansion of ε_0 .

On the other hand, the same infinitesimal variation computed through the left hand side

of Eq. (61) is just :

$$\begin{aligned}
& \varepsilon_0 \left(\mathbf{X}_{\hbar} + A_{\hbar}^{\mathbf{X}_{\hbar}} \right) - \mathcal{E} \varepsilon_0 \left(\mathbf{X}_{\hbar-d\hbar} + A_{\hbar-d\hbar}^{\mathbf{X}_{\hbar-d\hbar}} \right) \\
&= \mathcal{E} \sum_{n=0}^{\infty} \left[\int_{0 < \alpha_n < \dots < \alpha_1 < \hbar} e^{-\langle \cdot \rangle_{\alpha_1 \rightarrow \hbar}^S T_{\alpha_1}} \dots e^{-\langle \cdot \rangle_{\alpha_{n-1} \rightarrow \alpha_{n-2}}^S T_{\alpha_{n-1}}} e^{-\langle \cdot \rangle_{\alpha_n \rightarrow \alpha_{n-1}}^S T_{\alpha_n}} e^{-\langle \cdot \rangle_{0 \rightarrow \alpha_n}^S} \right] \varepsilon_0(\mathbf{X}_0) - \\
& \mathcal{E} e^{-\langle \cdot \rangle_{\hbar-d\hbar \rightarrow \hbar}^S} \sum_{n=0}^{\infty} \left[\int_{0 < \alpha_n < \dots < \alpha_1 < \hbar-d\hbar} e^{-\langle \cdot \rangle_{\alpha_1 \rightarrow \hbar-d\hbar}^S T_{\alpha_1}} \dots e^{-\langle \cdot \rangle_{\alpha_{n-1} \rightarrow \alpha_{n-2}}^S T_{\alpha_{n-1}}} e^{-\langle \cdot \rangle_{\alpha_n \rightarrow \alpha_{n-1}}^S T_{\alpha_n}} e^{-\langle \cdot \rangle_{0 \rightarrow \alpha_n}^S} \right] \varepsilon_0(\mathbf{X}_0) \\
&= [-\langle \cdot \rangle_{\hbar} + T_{\hbar}] \mathcal{E} \mathcal{T} \exp \left[\int_{0 < \alpha < \hbar-d\hbar} T_{\alpha} d\alpha \right] \varepsilon_0(\mathbf{X}) d\hbar \\
&+ \langle \cdot \rangle_{\hbar} \mathcal{E} \mathcal{T} \exp \left[\int_{0 < \alpha < \hbar-d\hbar} T_{\alpha} d\alpha \right] \varepsilon_0(\mathbf{X}) d\hbar \tag{62}
\end{aligned}$$

so that finally

$$\varepsilon_0 \left(\mathbf{X}_{\hbar} + A_{\hbar}^{\mathbf{X}_{\hbar}} \right) - \mathcal{E} \varepsilon_0 \left(\mathbf{X}_{\hbar-d\hbar} + A_{\hbar-d\hbar}^{\mathbf{X}_{\hbar-d\hbar}} \right) = T_{\hbar} \mathcal{E} \mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} T_{\alpha} d\alpha \right] \varepsilon_0(\mathbf{X}) d\hbar \tag{63}$$

at the lowest order in $d\hbar$. As a consequence, one is left ultimately with :

$$T_{\hbar} \mathcal{E} \mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} T_{\alpha} d\alpha \right] \varepsilon_0(\mathbf{X}) = \frac{\partial A_{\hbar}^{\mathbf{X}}}{\partial \hbar} \nabla_{\mathbf{X}} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}}) + \langle \cdot \rangle_{\hbar} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}}) \tag{64}$$

so that (recall that $\mathbf{X} \equiv \mathbf{X}_{\hbar}$) one can write

$$[-\langle \cdot \rangle_{\hbar} + T_{\hbar}] \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}}) = \frac{dA_{\hbar}^{\mathbf{X}}}{d\hbar} \nabla_{\mathbf{X}} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}}) \tag{65}$$

Now consider $\mathcal{P}_+[\mathcal{A}_{\alpha}^{\mathbf{R}}]$ and $\mathcal{P}_+[\mathcal{A}_{\alpha}^{\mathbf{P}}]$, with the following definitions for the "non-projected" Berry connections $\mathcal{A}_{\alpha}^{\mathbf{R}} = i[U_{\alpha} \nabla_{\mathbf{P}} U_{\alpha}^{\dagger}]$ and $\mathcal{A}_{\alpha}^{\mathbf{P}} = -i[U_{\alpha} \nabla_{\mathbf{R}} U_{\alpha}^{\dagger}]$. Define also for convenience $\mathcal{P}_+[\mathcal{A}_{\alpha}^{\mathbf{X}}] = (\mathcal{P}_+[\mathcal{A}_{\alpha}^{\mathbf{R}}], \mathcal{P}_+[\mathcal{A}_{\alpha}^{\mathbf{P}}])$. Our last equation Eq. (65) is thus equivalent to :

$$\begin{aligned}
& \frac{1}{2} \mathcal{P}_+[\mathcal{A}_{\hbar}^{\mathbf{X}}] \cdot \left[\nabla_{\mathbf{X}} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}}) + \nabla_{\mathbf{X}} (A_{\hbar}^{\mathbf{X}})^i(\mathbf{X}) \nabla_{\mathbf{X}_i} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}}) \right] + H.C. - \langle \cdot \rangle_{\hbar} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}}) \\
&= \frac{dA_{\hbar}^{\mathbf{X}}(\mathbf{X})}{d\hbar} \nabla_{\mathbf{X}_{\varepsilon_0 0}}(\mathbf{X} + A_{\hbar}^{\mathbf{X}})
\end{aligned}$$

where $H.C.$ stands for the Hermitic conjugate. The term $\nabla_{\mathbf{X}} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}})$ stems for the gradient of $\varepsilon_0(\mathbf{X})$ evaluated at $(\mathbf{X} + A_{\hbar}^{\mathbf{X}})$. As before $\nabla_{\mathbf{X}} (A_{\hbar}^{\mathbf{X}})^i(\mathbf{X}) \nabla_{\mathbf{X}_i} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}})$ has to be understood in the sense that $\nabla_{\mathbf{X}} (A_{\hbar}^{\mathbf{X}})^i$ replaces $\mathbf{X} + A_{\hbar}^{\mathbf{X}}$ at each place the gradient is acting on the series expansion of ε_0 . To be able to compare both sides of this equation, one has to symmetrize both expressions in the same way. As a consequence the left hand side

has to be rewritten in the same symmetrized form as the right hand side :

$$\begin{aligned} & \frac{1}{2} \mathcal{P}_+ [\mathcal{A}_\hbar^{\mathbf{X}}] \cdot \left[\nabla_{\mathbf{X}} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}}) + \nabla_{\mathbf{X}} (A_\hbar^{\mathbf{X}})^i (\mathbf{X}) \nabla_{\mathbf{X}_i} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}}) \right] + H.C. \\ &= \frac{1}{2} \left\{ \mathcal{P}_+ [\mathcal{A}_\hbar^{\mathbf{X}}] \cdot \left[\nabla_{\mathbf{X}} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}}) + \nabla_{\mathbf{X}} (A_\hbar^{\mathbf{X}})^i (\mathbf{X}) \nabla_{\mathbf{X}_i} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}}) \right] \right\}_S + H.C. + C(\mathbf{X}) \end{aligned}$$

This formula requires some explanation. In the left hand side, the multiplication by $\mathcal{P}_+ [\mathcal{A}_\alpha^{\mathbf{X}}]$ is performed half on the left and half on the right, as implied by definition of the translation operator. On the right hand side, all expressions are seen as symmetrized in a way that $\mathcal{P}_+ [\mathcal{A}_\alpha^{\mathbf{X}}]$ has been inserted at each place where the $\nabla_{\mathbf{X}}$ is acting (exactly as for $\frac{dA_\hbar^{\mathbf{X}}(\mathbf{X})}{d\hbar} \nabla_{\mathbf{X}} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}})$). The $\{\cdot\}_S$ is there to recall this full symmetrization in the variables. The term $C(\mathbf{X})$ is the correction due to this change of symmetrization while moving the $\mathcal{P}_+ [\mathcal{A}_\alpha^{\mathbf{X}}]$ inside the series expansion of ε_0 . By construction, it involves the powers of gradients of $\mathcal{P}_+ [\mathcal{A}_\alpha^{\mathbf{X}}]$ and ε_0 and is of order \hbar^2 . In practice, this term depends specifically on the problem at hand (that is on the form of the Hamiltonian) and can be computed order by order in \hbar . Since later on we will consider only the order \hbar^2 it will turn out that this term will be negligible. Let us just mention that, by construction, at the lowest order in \hbar , $C(\mathbf{X})$ is equal to $\langle \cdot \rangle_{\hbar} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}})$

We can now write the differential equation :

$$\begin{aligned} \frac{dA_\hbar^{\mathbf{X}}(\mathbf{X})}{d\hbar} \nabla_{\mathbf{X}} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}}) &= \frac{1}{2} \left\{ \mathcal{P}_+ [\mathcal{A}_\alpha^{\mathbf{X}}] \cdot \left[\nabla_{\mathbf{X}} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}}) + \nabla_{\mathbf{X}} (A_\hbar^{\mathbf{X}})^i (\mathbf{X}) \nabla_{\mathbf{X}_i} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}}) \right] \right\}_S \\ &+ H.C. + C(\mathbf{X}) - \langle \cdot \rangle_{\hbar} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}}) \end{aligned}$$

where all expressions are now symmetrized in the same way. Now, remark that, given our previous remarks, $C(\mathbf{X}) - \langle \cdot \rangle_{\hbar} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}})$ is of order \hbar^3 (since, as said before $C(\mathbf{X}) = \langle \cdot \rangle_{\hbar} \varepsilon_0 (\mathbf{X} + A_\hbar^{\mathbf{X}})$ at the lowest order, i.e. \hbar^2). It implies that, after integration, it will contribute only to the fourth order in \hbar to $A_\hbar^{\mathbf{X}}$. As a consequence, neglecting this term in first approximation (that will be indeed the case in our applications) we thus deduce that $A_\hbar^{\mathbf{X}}$ satisfies the following differential equation :

$$\mathcal{P}_+ [\mathcal{A}_\hbar^{\mathbf{X}}] + \frac{1}{2} (\mathcal{P}_+ [\mathcal{A}_\hbar^{\mathbf{X}}] \cdot \nabla_{\mathbf{X}}) A_\hbar^{\mathbf{X}} + H.C. = \frac{\partial A_\hbar^{\mathbf{X}}}{\partial \hbar} \quad (66)$$

whose solution can be written recursively as:

$$\begin{aligned} A_\hbar^{\mathbf{X}} &= \int_{0 < \alpha < \hbar} \left[\mathcal{P}_+ [\mathcal{A}_\alpha^{\mathbf{X}}] + \frac{1}{2} ((\mathcal{P}_+ [\mathcal{A}_\alpha^{\mathbf{X}}] \cdot \nabla_{\mathbf{X}}) A_\alpha^{\mathbf{X}} + H.C.) \right] d\alpha \\ &= \int_{0 < \alpha < \hbar} \mathcal{P}_+ [\mathcal{A}_\alpha^{\mathbf{X}}] d\alpha + \int_{0 < \alpha < \hbar} \frac{1}{2} \left[\left[\mathcal{P}_+ [\mathcal{A}_\alpha^{\mathbf{X}}] \cdot \nabla_{\mathbf{X}} \int_{0 < \alpha_1 < \alpha} \mathcal{P}_+ [\mathcal{A}_{\alpha_1}^{\mathbf{X}}] \right] + H.C. \right] d\alpha_1 d\alpha + \dots \end{aligned}$$

Let us stress that in the previous integrals, the variables \mathbf{X}_α in $A_\alpha^{\mathbf{R}}$ and $A_\alpha^{\mathbf{P}}$ have been replaced by \mathbf{X} and are thus constant with respect to the α integrations. The reason is that in the differential equation Eq. (66), \mathbf{X} is seen as constant, only \hbar is running. Actually, the equation involves only the partial derivative $\frac{\partial A_{\hbar}^{\mathbf{X}}}{\partial \hbar}$. Having now the solution for $A_{\hbar}^{\mathbf{X}}$ as a function of the Berry phases we can write our solution for the exponentiated action of the translation operator. Dividing $A_{\hbar}^{\mathbf{X}}$ in two components with respect to \mathbf{R} and \mathbf{P} , $A_{\hbar}^{\mathbf{X}} \equiv (A_{\mathbf{R}}, A_{\mathbf{P}})$, one has :

$$\mathcal{E}\mathcal{T} \exp \left[\int_{0 < \alpha < \hbar} e^{-\langle \cdot \rangle_{\alpha \rightarrow \hbar}^S} T_{\alpha} e^{-\langle \cdot \rangle_{\hbar \rightarrow \alpha}^S} d\alpha \right] \varepsilon_0(\mathbf{X}) = \varepsilon_0(\mathbf{x}) \quad (67)$$

with $\mathbf{x} = (\mathbf{r}, \mathbf{p})$ and \mathbf{r} and \mathbf{p} are new coordinate and momentum operators corrected by Berry connections terms in the following way :

$$\begin{aligned} \mathbf{r} &\equiv \mathbf{R} + A_{\mathbf{R}} \\ \mathbf{p} &\equiv \mathbf{P} + A_{\mathbf{P}} \end{aligned} \quad (68)$$

justifying the name translation operator for T_{α} . The inclusion of the corrections due to $C(\mathbf{X})$ can be performed in the following way. Shifting $A_{\hbar}^{\mathbf{X}}$ by a correction $A_{\hbar}^{\mathbf{X}} + \delta A_{\hbar}^{\mathbf{X}}$ with $A_{\hbar}^{\mathbf{X}}$ the solution previously found for Eq. (66), gives the following equation for $\delta A_{\hbar}^{\mathbf{X}}$:

$$\begin{aligned} &\frac{1}{2} \mathcal{P}_+ [A_{\hbar}^{\mathbf{X}}] \cdot \left[\nabla_{\mathbf{X}} (\delta A_{\hbar}^{\mathbf{X}})^i(\mathbf{X}) \nabla_{\mathbf{X}_i} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}} + \delta A_{\hbar}^{\mathbf{X}}) \right] + H.C. + C(\mathbf{X}) - \langle \cdot \rangle_{\hbar} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}}) \\ &= \frac{\partial \delta A_{\hbar}^{\mathbf{X}}(\mathbf{X})}{\partial \hbar} \nabla_{\mathbf{X}} \varepsilon_0(\mathbf{X} + A_{\hbar}^{\mathbf{X}} + \delta A_{\hbar}^{\mathbf{X}}) \end{aligned}$$

Note that $C(\mathbf{X})$ is computed using the derivatives of $A_{\hbar}^{\mathbf{X}} + \delta A_{\hbar}^{\mathbf{X}}$. Expanding $\delta A_{\hbar}^{\mathbf{X}}(\mathbf{X})$ as an \hbar series expansion of completely symmetrized function $\delta A_{\hbar}^{\mathbf{X}^{(n)}}(\mathbf{X})$ of \mathbf{R} and \mathbf{P} (starting with $n = 4$) allows, at least theoretically to find the $\delta A_{\hbar}^{\mathbf{X}^{(n)}}$ recursively as solving an equation of the kind :

$$\frac{\partial \delta A_{\hbar}^{\mathbf{X}^{(n+1)}}(\mathbf{X})}{\partial \hbar} = F\left(\mathbf{X}, \delta A_{\hbar}^{\mathbf{X}^{(n)}}(\mathbf{X}), \dots, \delta A_{\hbar}^{\mathbf{X}^{(n)}}(\mathbf{X})\right)$$

where $F\left(\mathbf{X}, \delta A_{\hbar}^{\mathbf{X}^{(n)}}(\mathbf{X}), \dots, \delta A_{\hbar}^{\mathbf{X}^{(n)}}(\mathbf{X})\right)$ is determined by replacing ε_0 , $A_{\hbar}^{\mathbf{X}}$, $\mathcal{P}_+ [A_{\hbar}^{\mathbf{X}}]$ by their expressions. As said before, in our practical considerations the term $\delta A_{\hbar}^{\mathbf{X}}(\mathbf{X})$ will always be neglected.

The variables $\mathbf{x} = \mathbf{X} + A_{\hbar}^{\mathbf{X}}$ but have defined seem thus to be the natural ones arising in the diagonalization process.

Note that at the lowest order we have :

$$\begin{aligned}\mathbf{r} &= \mathbf{R} + \hbar\mathcal{P}_+ [\mathcal{A}_0^{\mathbf{R}}] \equiv \mathbf{R} + A_0^{\mathbf{R}} \\ \mathbf{p} &= \mathbf{P} + \hbar\mathcal{P}_+ [\mathcal{A}_0^{\mathbf{P}}] \equiv \mathbf{P} + A_0^{\mathbf{P}}\end{aligned}\quad (69)$$

where $A_0^{\mathbf{R}}$ and $A_0^{\mathbf{P}}$ are the usual "semiclassical Berry connections" defined previously [5].

From equations Eq. (68) we readily deduce the following non trivial algebra

$$\begin{aligned}[r_i, r_j] &= i\hbar^2\Theta_{ij}^{rr} = i\hbar^2 (\nabla_{P_i}A_{R_j} - \nabla_{P_j}A_{R_i}) + \hbar^2 [A_{R_j}, A_{R_i}] \\ [p_i, p_j] &= i\hbar^2\Theta_{ij}^{pp} = -i\hbar^2 (\nabla_{R_i}A_{P_j} - \nabla_{R_j}A_{P_i}) + \hbar^2 [A_{P_i}, A_{P_j}] \\ [p_i, r_j] &= -i\hbar\delta_{ij} + i\hbar^2\Theta_{ij}^{pr} = -i\hbar\delta_{ij} - i\hbar^2 (\nabla_{R_i}A_{R_j} + \nabla_{P_j}A_{P_i}) + \hbar^2 [A_{P_i}, A_{R_j}]\end{aligned}\quad (70)$$

where the terms Θ_{ij} are the definitions of Berry curvatures. Of course these non trivial commutation relations also give new contributions to the equations of motion and thus lead to new phenomena [5] [7][11]. The commutation relations are valid to any order in \hbar , but in practice we can compute them as well as the energy $\varepsilon(\mathbf{X})$ in a series expansion in \hbar . Relations Eqs. (59) (67) will be very helpful when writing the explicit expression of $\varepsilon(\mathbf{X})$ in a series expansion in \hbar in the following section.

VI. SERIES EXPANSION IN \hbar

The exact expression Eq. (35), can now be expanded in a series expansion in \hbar . Note that we will always identify $\varepsilon_{\hbar}(\mathbf{R}_{\hbar}, \mathbf{P}_{\hbar})$ with $\varepsilon(\mathbf{R}, \mathbf{P})$.

A. First order in \hbar

At the first order we obviously get the following expression :

$$\begin{aligned}\varepsilon(\mathbf{X}) &= \varepsilon_0(\mathbf{X}) + \int_0^{\hbar} O_{\alpha} d\alpha \varepsilon_0(\mathbf{X}_{\alpha}) \\ &= \varepsilon_0(\mathbf{X}) + \hbar(T_{\hbar} + M_{\hbar})\varepsilon_0(\mathbf{X}) \\ &= \varepsilon_0(\mathbf{x}) + \frac{i\hbar}{2}\mathcal{P}_+ \left\{ [\varepsilon_0(\mathbf{X}), \mathcal{A}_0^{R_l}] \mathcal{A}_0^{P_l} - [\varepsilon_0(\mathbf{X}), \mathcal{A}_0^{P_l}] \mathcal{A}_0^{R_l} \right\} + O(\hbar^2)\end{aligned}\quad (71)$$

One then recover the formula first derived in ref. [5] where $x = (\mathbf{r}, \mathbf{p})$ are given by the expression Eq. (68).

As previously mentioned we only need to have $U_0(\mathbf{Y})$ and thus the zeroth order Berry phases to get the energy expansion at the first order. Now with the first order expansion for the diagonalization matrix :

$$\begin{aligned}
U(\mathbf{X}) &= U_0(\mathbf{X}) + \int_0^{\hbar} N_\alpha d\alpha U_0(\mathbf{X}) \equiv U_0(\mathbf{X}) + \hbar U_1(\mathbf{X}) U_0(\mathbf{X}) \\
&= \left(1 - \hbar [\cdot, \varepsilon_0]^{-1} \cdot \left[\mathcal{P}_- \left\{ \frac{1}{2} \left(\mathcal{A}_0^{R_l} \nabla_{R_l} \varepsilon_0(\mathbf{X}) + \nabla_{R_l} \varepsilon_0(\mathbf{X}) \mathcal{A}_0^{R_l} + \mathcal{A}_0^{P_l} \nabla_{P_l} \varepsilon_0(\mathbf{X}) + \nabla_{P_l} \varepsilon_0(\mathbf{X}) \mathcal{A}_0^{P_l} \right) \right\} \right. \right. \\
&\quad \left. \left. - \frac{i}{2} \left\{ [\varepsilon_0(\mathbf{X}), \mathcal{A}_0^{R_l}] \mathcal{A}_0^{P_l} - [\varepsilon_0(\mathbf{X}), \mathcal{A}_0^{P_l}] \mathcal{A}_0^{R_l} \right\} \right] - \frac{i}{4} [\mathcal{A}_0^{R_l}, \mathcal{A}_0^{P_l}] \right) U_0(\mathbf{X}) \quad (72)
\end{aligned}$$

B. Second order

At this order, we implicitly assume for convenience that all expressions are symmetrized in \mathbf{R} and \mathbf{P} , in such a way that for all expression depending on \mathbf{R} and \mathbf{P} , in the series expansion of this expression, all powers of the momentum are put half on the left and half on the right. Any other choice of symmetrization would be of course suitable.

To the second order, expanding the compact form of the energy operator leads to :

$$\varepsilon(\mathbf{X}) = \varepsilon_0(\mathbf{X}) + \int_0^{\hbar} O_\alpha d\alpha \varepsilon_0(\mathbf{X}_\alpha) + \int_0^{\hbar} O_{\alpha_1} \int_0^{\alpha_1} O_{\alpha_2} d\alpha_2 d\alpha_1 \varepsilon_0(\mathbf{X}_{\alpha_2}) - \frac{\hbar}{2} \langle \varepsilon_0(\mathbf{X}) \rangle$$

The last contribution is simply

$$-\frac{\hbar}{2} \langle \varepsilon_0(\mathbf{X}) \rangle = \frac{i}{4} \hbar \text{Asym} \{ \nabla_{P_l} \nabla_{R_l} \varepsilon_0(\mathbf{X}) \}$$

The reason for the $\frac{1}{2}$ factor is the following. Due to our choice of symmetrization, the bracket $\langle \cdot \rangle_\alpha$ induces an α factor. The integration over α yields then the $\frac{1}{2}$. The first contribution $\int_0^{\hbar} O_\alpha d\alpha \varepsilon_0(\mathbf{X})$ can be expanded as before as:

$$\begin{aligned}
&\int_0^{\hbar} O_\alpha d\alpha \varepsilon_0(\mathbf{X}) \\
&= \int_0^{\hbar} \mathcal{P}_+ \left\{ \frac{1}{2} \left(\mathcal{A}_\alpha^{R_l} \nabla_{R_l} \varepsilon_0(\mathbf{X}_\alpha) + \nabla_{R_l} \varepsilon_0(\mathbf{X}_\alpha) \mathcal{A}_\alpha^{R_l} + \mathcal{A}_\alpha^{P_l} \nabla_{P_l} \varepsilon_0(\mathbf{X}_\alpha) + \nabla_{P_l} \varepsilon_0(\mathbf{X}_\alpha) \mathcal{A}_\alpha^{P_l} \right) \right\} d\alpha \\
&+ \int_0^{\hbar} \mathcal{P}_+ \left\{ \frac{i}{4} [\varepsilon_0(\mathbf{X}_\alpha), \mathcal{A}_\alpha^{R_l}] \mathcal{A}_\alpha^{P_l} - \frac{i}{4} [\varepsilon_0(\mathbf{X}_\alpha), \mathcal{A}_\alpha^{P_l}] \mathcal{A}_\alpha^{R_l} + H.C. \right. \\
&\left. + \left[U_\alpha \left(\left(\frac{\partial}{\partial \alpha} + \langle \cdot \rangle \right) H(\mathbf{X}_\alpha) \right) U_\alpha^+ \right] \right\} d\alpha
\end{aligned}$$

where the Berry connections have to be expanded to the first order, whereas the second order contribution

$$\int_0^{\hbar} O_{\alpha_1} \int_0^{\alpha_1} O_{\alpha_2} d\alpha_2 d\alpha_1 \varepsilon_0(\mathbf{X}) - \frac{\hbar}{2} \langle \varepsilon_0(\mathbf{X}) \rangle$$

has to be expanded to the zeroth order in the Berry connections.

Notice that due to the integration process, the squared terms in $\mathcal{A}_0^{\mathbf{R}}$, $\mathcal{A}_0^{\mathbf{P}}$ as well as the first order terms in the Berry phase $\mathcal{A}_1^{\mathbf{R}}$ and $\mathcal{A}_1^{\mathbf{P}}$ get a $\frac{1}{2}$ factor. Defining the integrated Berry phases as :

$$\begin{aligned} \mathcal{A}^{R_l}(\mathbf{X}) &= \int_{\alpha}^{\hbar} \mathcal{A}_{\alpha}^{R_l} d\alpha \\ \mathcal{A}^{P_l}(\mathbf{X}) &= \int_{\alpha}^{\hbar} \mathcal{A}_{\alpha}^{P_l} d\alpha \end{aligned}$$

The contributions to the Hamiltonian can be recombined to yield :

$$\begin{aligned} \varepsilon(\mathbf{X}) &= \varepsilon_0(\mathbf{x}) + \frac{i}{2} \mathcal{P}_+ \left\{ \left[\varepsilon_0(\mathbf{x}), \hat{\mathcal{A}}^{R_l} \right] \hat{\mathcal{A}}^{P_l} - \left[\varepsilon_0(\mathbf{x}), \hat{\mathcal{A}}^{P_l} \right] \hat{\mathcal{A}}^{R_l} - \left[\varepsilon_0(\mathbf{x}), \left[\hat{\mathcal{A}}^{R_l}, \hat{\mathcal{A}}^{P_l} \right] \right] \right\} \\ &\quad - \frac{1}{8} \mathcal{P}_+ \left\{ \left[\left[\varepsilon_0(\mathbf{x}), \hat{\mathcal{A}}^{R_l} \right] \hat{\mathcal{A}}^{P_l} - \left[\varepsilon_0(\mathbf{x}), \hat{\mathcal{A}}^{P_l} \right] \hat{\mathcal{A}}^{R_l}, \hat{\mathcal{A}}^{R_l} \right] \hat{\mathcal{A}}^{P_l} \right\} \\ &\quad + \frac{1}{8} \mathcal{P}_+ \left\{ \left[\left[\varepsilon_0(\mathbf{x}), \hat{\mathcal{A}}^{R_l} \right] \hat{\mathcal{A}}^{P_l} - \left[\varepsilon_0(\mathbf{x}), \hat{\mathcal{A}}^{P_l} \right] \hat{\mathcal{A}}^{R_l}, \hat{\mathcal{A}}^{P_l} \right] \hat{\mathcal{A}}^{R_l} \right\} - \frac{\hbar}{2} \langle \varepsilon_0(\mathbf{x}) \rangle \end{aligned} \quad (73)$$

with :

$$\begin{aligned} \hat{\mathcal{A}}^{R_l} &= \frac{1}{2} \left[1 - \frac{1}{2} (\mathcal{P}_+ \mathcal{A}_0^{\mathbf{X}} \cdot \nabla_{\mathbf{x}}) \right] \mathcal{A}^{R_l}(\mathbf{x}) + H.C. \\ \hat{\mathcal{A}}^{P_l} &= \frac{1}{2} \left[1 - \frac{1}{2} (\mathcal{P}_+ \mathcal{A}_0^{\mathbf{X}} \cdot \nabla_{\mathbf{x}}) \right] \mathcal{A}^{P_l}(\mathbf{x}) + H.C. \end{aligned}$$

Note that in the all Hamiltonian, including the Berry connections \mathcal{A}^{R_l} and \mathcal{A}^{P_l} , we have replaced the operators (\mathbf{R}, \mathbf{P}) by $\mathbf{x} = (\mathbf{r}, \mathbf{p})$ at each order of the expansion.

We can now recall our previous definition of the variables $\mathbf{x} = \mathbf{X} + \mathcal{A}^{\mathbf{X}}$:

$$\begin{aligned}
\mathbf{r} &= \mathbf{R} + \int_0^{\hbar} \mathcal{P}_+ \mathcal{A}_{\alpha_1}^{\mathbf{R}}(\mathbf{R}, \mathbf{P}) d\alpha_1 + \int_{0 < \alpha < \hbar} \frac{1}{2} \left[\left[\mathcal{P}_+ [\mathcal{A}_{\alpha}^{\mathbf{X}}] \cdot \nabla_{\mathbf{X}} \int_{0 < \alpha_1 < \alpha} \mathcal{P}_+ [\mathcal{A}_{\alpha_1}^{\mathbf{R}}] + H.C. \right] d\alpha_1 d\alpha \right. \\
&= \mathbf{R} + \mathcal{P}_+ \mathcal{A}^{\mathbf{R}}(\mathbf{X}) + \frac{\hbar^2}{4} \left((\mathcal{P}_+ \mathcal{A}_0^{\mathbf{R}} \cdot \nabla_{\mathbf{R}}) \mathcal{P}_+ \mathcal{A}_0^{\mathbf{R}} + (\mathcal{P}_+ \mathcal{A}_0^{\mathbf{P}} \cdot \nabla_{\mathbf{P}}) \mathcal{P}_+ \mathcal{A}_0^{\mathbf{R}} + H.C. \right) \\
&\equiv \mathbf{R} + \hbar \mathcal{A}_0^{\mathbf{R}} + \frac{\hbar^2}{2} \mathcal{A}_1^{\mathbf{R}} \\
\mathbf{p} &= \mathbf{P} + \int_0^{\hbar} \mathcal{P}_+ \mathcal{A}_{\alpha_1}^{\mathbf{P}}(\mathbf{R}, \mathbf{P}) d\alpha_1 + \int_{0 < \alpha < \hbar} \frac{1}{2} \left[\left[\mathcal{P}_+ [\mathcal{A}_{\alpha}^{\mathbf{X}}] \cdot \nabla_{\mathbf{X}} \int_{0 < \alpha_1 < \alpha} \mathcal{P}_+ [\mathcal{A}_{\alpha_1}^{\mathbf{P}}] + H.C. \right] d\alpha_1 d\alpha \right. \\
&= \mathbf{P} + \mathcal{P}_+ \mathcal{A}^{\mathbf{P}} + \frac{\hbar^2}{4} \left(\frac{1}{2} (\mathcal{P}_+ \mathcal{A}_0^{\mathbf{R}} \cdot \nabla_{\mathbf{R}}) \mathcal{P}_+ \mathcal{A}_0^{\mathbf{P}} + (\mathcal{P}_+ \mathcal{A}_0^{\mathbf{P}} \cdot \nabla_{\mathbf{P}}) \mathcal{P}_+ \mathcal{A}_0^{\mathbf{P}} + H.C. \right) \\
&\equiv \mathbf{P} + \hbar \mathcal{A}_0^{\mathbf{P}} + \frac{\hbar^2}{2} \mathcal{A}_1^{\mathbf{P}} \tag{74}
\end{aligned}$$

which of course satisfy an algebra Eq. (70).

We end up this section by giving some expanded formula for the Berry phases $\mathcal{A}_{\alpha}^{\mathbf{X}}$ intervening in the definition of the diagonalized Hamiltonian and the dynamics variables. At this order, we only require the matrix U at the first order as computed before : $U(\mathbf{X}) \equiv U_{\hbar}(\mathbf{X}) = U_0(\mathbf{X}) + \hbar U_1(\mathbf{X}) U_0(\mathbf{X})$ where $U_1(\mathbf{X})$ has been given before

$$\begin{aligned}
U_1(\mathbf{X}) &= [., \varepsilon_0]^{-1} \cdot \left[\mathcal{P}_- \left\{ \frac{1}{2} \left(\mathcal{A}_0^{R_l} \nabla_{R_l} \varepsilon_0(\mathbf{X}) + \nabla_{R_l} \varepsilon_0(\mathbf{X}) \mathcal{A}_0^{R_l} + \mathcal{A}_0^{P_l} \nabla_{P_l} \varepsilon_0(\mathbf{X}) + \nabla_{P_l} \varepsilon_0(\mathbf{X}) \mathcal{A}_0^{P_l} \right) \right\} \right. \\
&\quad \left. - \frac{i}{2} \left\{ [\varepsilon_0(\mathbf{X}), \mathcal{A}_0^{R_l}] \mathcal{A}_0^{P_l} - [\varepsilon_0(\mathbf{X}), \mathcal{A}_0^{P_l}] \mathcal{A}_0^{R_l} \right\} \right] - \frac{i}{4} [\mathcal{A}_0^{R_l}, \mathcal{A}_0^{P_l}] \tag{75}
\end{aligned}$$

At the same order the (non-diagonal) Berry connections $\mathcal{A}_{\alpha}^{\mathbf{X}} = (\mathcal{A}_{\alpha}^{\mathbf{R}}, \mathcal{A}_{\alpha}^{\mathbf{P}})$ are again given by

$$\begin{aligned}
\mathcal{A}_{\alpha}^{\mathbf{R}}(\mathbf{X}_{\alpha}) &= i [U_{\alpha}(\mathbf{X}_{\alpha}) \nabla_{\mathbf{P}} U_{\alpha}^+(\mathbf{X}_{\alpha})] = \frac{1}{\alpha} U_{\alpha}(\mathbf{X}_{\alpha}) \mathbf{R}_{\alpha} U_{\alpha}^+(\mathbf{X}_{\alpha}) \tag{76} \\
\text{and } \mathcal{A}_{\alpha}^{\mathbf{P}}(\mathbf{X}_{\alpha}) &= -i [U_{\alpha}(\mathbf{X}_{\alpha}) \nabla_{\mathbf{R}} U_{\alpha}^+(\mathbf{X}_{\alpha})] = \frac{1}{\alpha} U_{\alpha}(\mathbf{X}_{\alpha}) \mathbf{P}_{\alpha} U_{\alpha}^+(\mathbf{X}_{\alpha})
\end{aligned}$$

where now $U_{\alpha}(\mathbf{X}_{\alpha})$ is the transformation to the first order in α i.e. $U_0(\mathbf{X}_{\alpha}) + \alpha U_1(\mathbf{X}_{\alpha})$, in which \mathbf{X} is replaced by the running operator \mathbf{X}_{α} . Using the hermiticity of $\mathcal{A}_{\alpha}^{\mathbf{X}}$, so that one has $\mathcal{A}_{\alpha}^{\mathbf{X}} = \mathcal{A}_{\alpha}^{\mathbf{X}} + (\mathcal{A}_{\alpha}^{\mathbf{X}})^{\dagger}$ we can expand $\mathcal{A}^{\mathbf{X}}$ as :

$$\mathcal{A}_{\alpha}^{\mathbf{X}} = \left(\frac{1}{2} ([1 + \alpha U_1(\mathbf{X}_{\alpha})] U_0(\mathbf{X}_{\alpha})) \frac{\mathbf{X}_{\alpha}}{\alpha} (U_0^+(\mathbf{X}_{\alpha}) [1 + \alpha U_1^+(\mathbf{X}_{\alpha})]) + \frac{1}{2} H.C. \right) - \frac{\mathbf{X}_{\alpha}}{\alpha}$$

(the $\frac{1}{\alpha}$ factor reminds that in our definition of $\mathcal{A}_\alpha^{\mathbf{X}}$ the Gradient with respect to \mathbf{X}_α is normalized, i.e. divided by α). After some recombination, the previous expression can be written in a more convenient form :

$$\begin{aligned} \mathcal{A}_\alpha^{\mathbf{X}} &= \frac{1}{2\alpha} U_0(\mathbf{X}_\alpha) [\mathbf{X}_\alpha, U_0^+(\mathbf{X}_\alpha)] + H.C. \\ &+ \frac{1}{2} [U_1(\mathbf{X}_\alpha) [U_0(\mathbf{X}_\alpha) \mathbf{X}_\alpha U_0^+(\mathbf{X}_\alpha)] + [U_0(\mathbf{X}_\alpha) \mathbf{X}_\alpha U_0^+(\mathbf{X}_\alpha)] U_1^+(\mathbf{X}_\alpha)] + H.C. - \mathbf{X}_\alpha \end{aligned}$$

using now the fact that at the lowest order $U_0(\mathbf{X}_\alpha) [\frac{\mathbf{X}_\alpha}{\alpha}, U_0^+(\mathbf{X}_\alpha)] = \mathcal{A}_0^{\mathbf{X}}(\mathbf{X}_\alpha)$, one has :

$$\begin{aligned} \mathcal{A}_\alpha^{\mathbf{X}} &= \frac{1}{2\alpha} [U_0(\mathbf{X}_\alpha) [\mathbf{X}_\alpha, U_0^+(\mathbf{X}_\alpha)] + H.C.] \\ &+ \frac{1}{2} [U_1(\mathbf{X}_\alpha) \mathcal{A}_0^{\mathbf{X}}(\mathbf{X}_\alpha) + \mathcal{A}_0^{\mathbf{X}}(\mathbf{X}_\alpha) U_1^+(\mathbf{X}_\alpha) + [\mathbf{X}_\alpha, U_1^+(\mathbf{X}_\alpha)] + H.C.] \end{aligned}$$

decomposing $U_1^+(\mathbf{X}_\alpha)$ into hermitian and antihermitian part we are thus led to :

$$\begin{aligned} \mathcal{A}_\alpha^{\mathbf{X}} &= \frac{1}{2\alpha} (U_0(\mathbf{X}_\alpha) [\mathbf{X}_\alpha, U_0^+(\mathbf{X}_\alpha)] + H.C.) + [\mathbf{X}_\alpha + \mathcal{A}_0^{\mathbf{X}}, ah(U_1^+(\mathbf{X}_\alpha))] \\ &+ (\mathcal{A}_0^{\mathbf{X}} H(U_1^+(\mathbf{X}_\alpha)) + h(U_1^+(\mathbf{X}_\alpha)) \mathcal{A}_0^{\mathbf{X}}) \end{aligned}$$

where $ah(Z)$ and $h(Z)$ denote the antihermitian and hermitian part of an operator Z respectively. Now using Eq. (38), we are led to the following expression :

$$\mathcal{A}_\alpha^{\mathbf{X}} = \mathcal{A}_0^{\mathbf{X}} \left(\mathbf{R} + \frac{\alpha}{2} \mathcal{A}_0^{R_i}, \mathbf{P} + \frac{\alpha}{2} \mathcal{A}_0^{P_i} \right) - [B, \mathbf{X}_\alpha + \mathcal{A}_0^{\mathbf{X}}] \quad (77)$$

where we introduced the notations

$$\mathcal{A}_0^{\mathbf{X}} \left(\mathbf{R} + \frac{\alpha}{2} \mathcal{A}_0^{R_i}, \mathbf{P} + \frac{\alpha}{2} \mathcal{A}_0^{P_i} \right) \equiv \mathcal{A}_0^{\mathbf{X}} + \frac{\alpha}{4} \left\{ \mathcal{A}_0^{R_i} \nabla_{R_i} \mathcal{A}_0^{\mathbf{X}} + \mathcal{A}_0^{P_i} \nabla_{P_i} \mathcal{A}_0^{\mathbf{X}} + H.C. \right\} \quad (78)$$

and B is given by the following relation

$$\begin{aligned} B &= [., \varepsilon_0]^{-1} \cdot \left(\mathcal{P}_- \left\{ \frac{1}{2} \mathcal{A}_0^{R_i} \nabla_{R_i} \varepsilon_0(\mathbf{X}) + \frac{1}{2} \mathcal{A}_0^{P_i} \nabla_{P_i} \varepsilon_0(\mathbf{X}) + H.C. \right\} \right. \\ &\quad \left. - \frac{i}{4} \left\{ [\varepsilon_0(\mathbf{X}), \mathcal{A}_0^{R_i}] \mathcal{A}_0^{P_i} - [\varepsilon_0(\mathbf{X}), \mathcal{A}_0^{P_i}] \mathcal{A}_0^{R_i} + H.C. \right\} \right) \\ &= [., \varepsilon_0]^{-1} \cdot \left(\mathcal{P}_- \left\{ \frac{1}{2} \mathcal{A}_0^{R_i} \nabla_{R_i} \varepsilon_0(\mathbf{X}) + \frac{1}{2} \mathcal{A}_0^{P_i} \nabla_{P_i} \varepsilon_0(\mathbf{X}) + H.C. \right\} \right) \\ &\quad - \frac{i}{4} \left\{ \mathcal{P}_- \mathcal{A}_0^{R_i} \mathcal{P}_+ \mathcal{A}_0^{P_i} - \mathcal{P}_- \mathcal{A}_0^{P_i} \mathcal{P}_+ \mathcal{A}_0^{R_i} + H.C. \right\} \end{aligned} \quad (79)$$

VII. PHYSICAL APPLICATION: THE DIRAC ELECTRON IN AN ELECTRIC FIELD

To illustrate our general theory we consider the case of a Dirac electron in an external electric field. We will obtain the block diagonal Hamiltonian to the second order in \hbar and will compare with the FW transformation obtained in [8]. Note that contrary to the FW which is not an expansion in \hbar , the new method is valid for strong external fields (actually a FW transformation expanded into a power series in \hbar was also recently proposed [10]).

Let consider the following Dirac Hamiltonian ($c = 1$)

$$H_1 = \alpha \cdot \mathbf{P} + \beta m + eV(\mathbf{R})$$

At the first order in \hbar we just need the zero order transformation

$$U_0 = \frac{E + m + \beta \alpha \mathbf{P}}{\sqrt{2E(E + m)}}$$

where $E = \sqrt{\mathbf{P}^2 + m^2}$. In this case we have (with $\Sigma = \mathbf{1} \otimes \sigma$)

$$\begin{aligned} \mathcal{A}_0^R &= U_0 \mathbf{R} U_0^{-1} = i\hbar \frac{-\beta \alpha \cdot \mathbf{P} \mathbf{P} + E(E + m)\beta \alpha - i\mathbf{E} \mathbf{P} \times \Sigma}{2E^2(E + m)} \\ \mathcal{A}_0^P &= U_0^i \mathbf{P} U_0^{-1} = 0 \end{aligned}$$

which leads to the Berry connections (we use the block diagonal projector \mathcal{P}_+):

$$\begin{aligned} A_0^{\mathbf{R}} &= \frac{\mathbf{P} \times \Sigma}{2E(E + m)} \\ A_0^{\mathbf{P}} &= 0 \end{aligned}$$

For the second order we need the first order unitary matrix we can be deduce directly from Eq. (72)

$$U_1 = \left[\left(1 + \frac{\beta e}{2E} \mathcal{P}_- (\mathcal{A}_0^R) \cdot \nabla_R V \right) \right] U_0$$

From U_1 and using Eq. (74) we can compute the projected Berry connections to the second order in powers of \hbar and thus write the dynamical operators as

$$\mathbf{r} = \mathbf{R} + \hbar A_0^{\mathbf{R}} + \frac{\hbar^2}{2} A_1^{\mathbf{R}} = \hbar \frac{\mathbf{P} \times \Sigma}{2E(E + m)} + \frac{\hbar^2 e}{2} \beta \frac{E^2 \nabla_R V - [\mathbf{P} \cdot \nabla_R V] \mathbf{P}}{4E^5}$$

$$\mathbf{p} = \mathbf{P}$$

Now using expression Eq. (73) we arrive at the following expression for the diagonal representation of the energy operator

$$\begin{aligned}\varepsilon &= \beta E(\mathbf{r}, \mathbf{p}) - \beta E(\mathbf{r}, \mathbf{p}) \Theta_{ij}^{pr}(\mathbf{r}, \mathbf{p}) + V(\mathbf{r}) \\ &= \beta \sqrt{\mathbf{p}^2 + m^2} + \frac{\hbar^2 e}{2} \beta \nabla_{\mathbf{r}} \cdot \frac{E^2 \nabla_{\mathbf{r}} V - [\mathbf{p} \cdot \nabla_{\mathbf{r}} V] \mathbf{p}}{4E^4} + eV(\mathbf{r})\end{aligned}\quad (80)$$

Here we have made the choice of fully symmetrizing in \mathbf{r} and \mathbf{p} , that is to weight equally all permutations in \mathbf{r} and \mathbf{p} in the series expansions of our expressions. One can now check that developing the variables \mathbf{r} as a function of the canonical variables \mathbf{R} and \mathbf{P} yields the same expression for the Hamiltonian as in [8].

Eq. (80) being fully relativistic, to compare with the usual FW [1] approach we consider the non-relativistic limit and expand our results to second order in $\frac{1}{mc}$. We readily obtain the well known diagonal representation of the positive energy (expressed in coordinates R and P) [1] :

$$\varepsilon = \frac{\mathbf{P}^2}{2m} - \frac{\mathbf{P}^4}{8m^3 c^2} + eV(\mathbf{R}) + \frac{e\hbar}{4m^2 c^2} \sigma \cdot (\nabla_{\mathbf{R}} V \times \mathbf{p}) + \frac{e\hbar^2}{8m^2 c^2} \nabla^2 V \quad (81)$$

with σ the Pauli matrices. Note that it is the term order \hbar^2 in Eq. (80) which, in the non relativistic limit, leads to the Darwin term.

VIII. CONCLUSION

In this paper, we presented a new diagonalization method for a generic matrix valued Hamiltonian which leads to a diagonal representation where the operator energy takes an elegant and very compact form. This approach requires the introduction of some new mathematical objects like non-commuting operators evolving with the Planck constant promoted as a running variable and thus reveals a very rich mathematical structure reminiscent of the stochastic calculus. It also shows once more the very important role played by Berry phases in these systems as the energy operator is written in terms of covariant dynamical operators containing Berry connections and satisfying a non-commutative algebra.

It was also found that the diagonal representation of the energy is solution of a differential equation in \hbar presented previously in [5] and which could only be solved recursively in a series expansion in \hbar . Actually the formal exact solution presented here can also be written explicitly as a series expansion in \hbar , but it appears that the derivation of the coefficients

of this expansion is now much more easier. Indeed we could give the expression of the energy and the dynamical variables to the second order in \hbar for a generic matrix valued Hamiltonian. We then applied this method to the simple case of a Dirac electron in an external electric field and recover the usual Pauli-Hamiltonian in the non relativistic limit including the Darwin term which is a \hbar^2 contribution. This is obviously a good check for the validity of the proposed method.

This approach is powerful since it leads to the diagonal expression for the energy operator of any kind of quantum system interacting with external fields (which can be strong) and turns out to be more tractable for practical applications than any other previous methods. We leave for subsequent work its application to more complicated systems in condensed matter or relativistic particle physics such as Bloch and Dirac electrons in interaction.

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