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Perturbation theory and control in classical or quantum mechanics by an inversion formula

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Abstract
We consider a perturbation of an 'integrable' Hamiltonian and give an expression for the canonical or unitary transformation which 'simplifies' this perturbed system. The problem is to invert a functional defined on the Lie-algebra of observables. We give a bound for the perturbation in order to solve this inversion, and apply this result to a particular case of the control theory, as a first example, and to the 'quantum adiabatic transformation', as another example.

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1. Algebraic framework

We start with a vector space $\mathcal{A}$ which we call the space of observables. We will apply this theory to two main examples: classical mechanics, in which $\mathcal{A}$ is an (Abelian) algebra of functions defined on the phase space, and quantum mechanics, in which $\mathcal{A}$ is an algebra of operators on some Hilbert space. By taking a basis of projectors in $\mathcal{A}$ we will consider these operators as a collection of 'matrix elements', i.e. also as functions.

We assume that $\mathcal{A}$ is endowed with a Lie structure, i.e. there exists a mapping (called the 'bracket') from $\mathcal{A}$ into the space $\mathcal{L}(\mathcal{A})$ of linear operators of $\mathcal{A}$:

$$\{\cdots\} : \mathcal{A} \to \mathcal{L}(\mathcal{A}) \qquad V \mapsto \{V\} \tag{1.1}$$

which satisfies

$$\forall V, W \in \mathcal{A} \qquad \{V\}W = -\{W\}V \tag{1.2}$$

$$\text{and} \quad \{\{V\}W\} = \{V\}\{W\} - \{W\}\{V\} \tag{1.3}$$

so that the 'bracket' mapping is linear in its argument. It is also antisymmetric by (1.2) and satisfies the Jacobi identity (1.3). We consider a fixed element H of $\mathcal{A}$ which we call a Hamiltonian. The 'motion' or the 'flow' generated by H is the 1-parameter group:

$$\forall t \in \mathbb{R} \qquad \mathrm{e}^{t\{H\}} : \mathcal{A} \to \mathcal{A}. \tag{1.4}$$

The exponential of $\{H\}$ is defined by the usual power series:

$$\mathrm{e}^{t\{H\}} := \sum_{n=0}^{\infty} \frac{\{t \cdot H\}^n}{n!} \tag{1.5}$$

and is an automorphism of the Lie structure, as is proved in (5.4).

Actually what is needed is only that $\{H\}$ is a linear operator 'affiliated' to $\mathcal{A}$ which means that (1.4) is still valid: $\{H\}$ is assumed to be the generator of a 1-parameter group of automorphisms, even if H is not an element of $\mathcal{A}$.

Our problem is to find a relation (for instance a 'conjugation') between the flow generated by a perturbation of H, denoted as $H + V$ for some $V \in \mathcal{A}$, and the flow of H. Poincaré called this problem 'the main problem of the dynamics'. Of course, this relation would be mainly useful if we have some information on the flow of the unperturbed Hamiltonian H.

We will study two types of problems.

Problem 1. *If we are permitted to modify the perturbed Hamiltonian $H + V$ by adding a 'small' term $f(V)$, we will try to find a relation between the flow of $H + V + f(V)$ and the flow of H. Then the term $f(V)$ will be called a 'control' term.*

Of course we want to exclude the trivial solution $f(V) = -V$ so we have added the supplementary condition on this term to be 'small', for instance quadratic in V. So we stabilize the perturbed Hamiltonian by a small control term, well adapted to the problem.

Problem 2. *If we are not permitted to modify the perturbed Hamiltonian (for instance if we investigate the motion of some planet), or we do not want to modify it, then we will try to find a good 'change of coordinates', i.e. an automorphism of $\mathcal{A}$, which connects the two flows: perturbed and unperturbed.*

We will search this automorphism under the form of some exponential $\mathrm{e}^{\{\Gamma W\}}$ where we define Γ as follows.

Let us make an important assumption on H, which will be satisfied when H is 'integrable', as seen below.

Hypothesis 1. *We assume that there exists a linear operator $\Gamma : \mathcal{A} \to \mathcal{A}$ such that*

$$\{H\}^2\Gamma = \{H\} \tag{1.6}$$

and then we build two other operators $\mathcal{N}$ and $\mathcal{R}$ by

$$\mathcal{N} := \{H\}\Gamma \qquad \mathcal{R} := \mathbf{1} - \mathcal{N}. \tag{1.7}$$

Hence Γ is a pseudo-inverse of $\{H\}$: let us recall that it is impossible to find a strict inverse of $\{H\}$ since it always has a non-trivial kernel (for instance $\{H\}H = 0$).

Any element V of $\mathcal{A}$ such that $\{H\}V = 0$ is constant under the flow of H:

$$\mathrm{e}^{t\{H\}}V = V. \tag{1.8}$$

So the vector space $\mathrm{Ker}\{H\}$ is called the set of 'constants of motion'. It is a sub-Lie-algebra of $\mathcal{A}$ since (using (1.3)):

$$\text{if} \quad \{H\}V = \{H\}W = 0 \qquad \text{then} \quad \{H\}\{V\}W = \{V\}\{H\}W + \{\{H\}V\}W = 0 + 0. \tag{1.9}$$

Let us also note that (1.6) can be rewritten as

$$\{H\}\mathcal{R} = 0 \tag{1.10}$$

which means that the range Rg$\mathcal{R}$ of the operator $\mathcal{R}$ is included in Ker$\{H\}$. The $\mathbf{1}$ is the identity in the algebra $\mathcal{L}(\mathcal{A})$ of endomorphisms of $\mathcal{A}$. The notation $\mathcal{R}$ designates the 'resonant part' and the notation $\mathcal{N}$ the 'non-resonant part'. Let us recall that $\{H\}, \mathcal{R}, \mathcal{N}, \Gamma$ are elements of $\mathcal{L}(\mathcal{A})$.

1.1. Application 1: classical mechanics

$\mathcal{A}$ is the algebra of $\mathcal{C}^\infty$ real-valued functions of (p, q) in some domain of $\mathbb{R}^{2L}$ or functions of $(A, \theta) \in \mathcal{H} \times \mathbb{T}^L$ for some domain $\mathcal{H}$ of $\mathbb{R}^L$ (L is the number of degrees of freedom). It is more convenient to make a Fourier transformation in θ, and so $\mathcal{A}$ may be taken as the algebra of functions V of $(A, \Delta) \in \mathcal{H} \times \mathbb{Z}^L$ into $\mathbb{C}$, such that $V(A, -\Delta) = V(A, \Delta)^*$. The Lie structure is the Poisson bracket. We can compute the operators $\mathcal{R}$ and Γ when H is 'integrable', i.e. when H is a function of the 'action variables' only, and does not depend on the 'angle variables':

$$H(A, \theta) = h(A) \tag{1.11}$$

or after a Fourier transformation in θ:

$$H(A, \Delta) = h(A) \cdot \delta_0(\Delta) \tag{1.12}$$

with δ being the Kronecker symbol. Indeed, let us denote by $\omega(A)$ the derivative (gradient) of h with respect to A. Then the bracket $\{H\} = \omega(A) \cdot \partial_\theta$ and so, after a Fourier transformation, it is given by

$$(\{H\}V)(A, \Delta) = \mathbf{i} \cdot (\omega(A) \cdot \Delta) \cdot V(A, \Delta) \tag{1.13}$$

on any element V of $\mathcal{A}$. We have used a scalar product $\omega(A) \cdot \Delta$, so that the operator $\mathcal{R}$ is

$$(\mathcal{R}V)(A, \Delta) = V(A, \Delta) \cdot \chi(\omega(A) \cdot \Delta = 0) \tag{1.14}$$

where we introduce a characteristic function $\chi(\omega(A) \cdot \Delta = 0)$ which is 1 when $\omega(A) \cdot \Delta = 0$ and is 0 otherwise. We could have written it as $\chi(\Delta \in \omega(A)^\perp \cap \mathbb{Z}^L)$. Similarly

$$(\mathcal{N}V)(A, \Delta) = V(A, \Delta) \cdot \chi(\omega(A) \cdot \Delta \neq 0). \tag{1.15}$$

Then the action of the operator Γ is given by

$$(\Gamma V)(A, \Delta) := \frac{\chi(\omega(A) \cdot \Delta \neq 0)}{\mathbf{i} \cdot (\omega(A) \cdot \Delta)} \cdot V(A, \Delta). \tag{1.16}$$

See section 6 for more details and section 7 where we introduce a norm.

1.2. Application 2: quantum mechanics

$\mathcal{A}$ is the algebra of operators on some separable Hilbert space. More precisely we assume that $\mathcal{A}$ has a unity (denoted by $\mathbf{1}$), and we consider a maximal family of mutually orthogonal projectors $P_A \in \mathcal{A}$ where A varies in some countable set $\mathcal{H}$. That means

$$\sum_{A \in \mathcal{H}} P_A = \mathbf{1} \tag{1.17}$$

$$\forall A, A' \in \mathcal{H} \qquad P_A \cdot P_{A'} = P_A \cdot \delta_{A,A'} \tag{1.18}$$

with δ being the Kronecker symbol. $\mathcal{R}$ and Γ are easily written when H can be 'diagonalized'. We choose H as follows:

$$H = \sum_{A\in\mathcal{H}} h(A)\cdot P_A \tag{1.19}$$

for some function $h : \mathcal{H} \to \mathbb{R}$. So $h(\mathcal{H})$ is the spectrum of H and the projectors are the spectral projectors of H. We can choose h to be the identity function by taking $\mathcal{H}$ as the spectrum of H.

Remark 1. Reciprocally, if $\mathcal{A}$ were endowed with an involution ('*', named 'adjonction') we could have started by giving a self-adjoint operator H with pure point spectrum, and then we could have taken the spectral projectors of H, as the family of projectors P_A.

Then any element V of $\mathcal{A}$ can be written as

$$V = \sum_{A,A'\in\mathcal{H}} V_{A,A'} \qquad \text{where} \qquad V_{A,A'} := P_A V P_{A'}. \tag{1.20}$$

It is convenient to introduce the set

$$\mathcal{G}(A) := \mathcal{H} - A := \{A' - A\}_{A'\in\mathcal{H}} \tag{1.21}$$

and to define

$$\forall A \in \mathcal{H} \qquad \forall \Delta \in \mathcal{G}(A) \qquad V(A,\Delta) := P_{A+\Delta} V P_A \tag{1.22}$$

i.e. $V(A,\Delta) = V_{A+\Delta,A}$. For instance

$$H(A,\Delta) = h(A)\cdot\delta_0(\Delta)\cdot P_A \tag{1.23}$$

which is similar to (1.12).

The Lie structure is given by the commutator. Actually we multiply it by $\mathbf{i}/\hbar$ for some constant $\hbar$ which has the correct dimensionality, so that $t\{H\}$ has no dimension

$$\forall V, W \in \mathcal{A} \qquad \{V\}(W) := \mathbf{i}(V\cdot W - W\cdot V)/\hbar. \tag{1.24}$$

Let us now turn to the action of the operators $\{H\}$ and $\mathcal{R}$

$$(\{H\}V)(A,\Delta) = \mathbf{i}\cdot(h(A+\Delta) - h(A))\cdot V(A,\Delta)/\hbar. \tag{1.25}$$

Hence

$$(\mathcal{R}V)(A,\Delta) = V(A,\Delta)\cdot\chi(h(A+\Delta) = h(A)) \tag{1.26}$$

and

$$(\mathcal{N}V)(A,\Delta) = V(A,\Delta)\cdot\chi(h(A+\Delta) \neq h(A)). \tag{1.27}$$

Then the operator Γ is given by

$$(\Gamma V)(A,\Delta) = \frac{\hbar\cdot\chi(h(A+\Delta)\neq h(A))}{\mathbf{i}\cdot(h(A+\Delta) - h(A))}\cdot V(A,\Delta). \tag{1.28}$$

See section 6 for more details and section 7 where we introduce a norm.

Remark. We have introduced an arbitrary constant $\hbar$ in (1.24). But we will note that the operator Γ always appears with a bracket around it and the pair $\{\Gamma\ldots\}$ is independent of $\hbar$, since this constant is in the numerator and the denominator.

2. The main theorems

Theorem 1. *The control problem (problem 1) is solved by an explicit formula. Let us first define the functions* F and $f : \mathcal{A} \to \mathcal{A}$ *by*

$$F(V) := \mathrm{e}^{-\{\Gamma V\}}\mathcal{R}V + \frac{1-\mathrm{e}^{-\{\Gamma V\}}}{\{\Gamma V\}}\mathcal{N}V \tag{2.1}$$

$$f(V) := F(V) - V. \tag{2.2}$$

Then we have

$$\forall t \in \mathbb{R} \qquad \mathrm{e}^{t\{H+V+f(V)\}} = \mathrm{e}^{-\{\Gamma V\}} \cdot \mathrm{e}^{t\{H\}} \cdot \mathrm{e}^{t\{\mathcal{R}V\}} \cdot \mathrm{e}^{\{\Gamma V\}}. \tag{2.3}$$

The meaning of the second term in (2.1) is as follows:

$$\frac{1-\mathrm{e}^{-\{\Gamma V\}}}{\{\Gamma V\}} = \sum_{n\in\mathbb{N}} \frac{\{-\Gamma V\}^n}{n+1!} = \int_0^1 \mathrm{d}s \cdot \mathrm{e}^{-s\{\Gamma V\}}. \tag{2.4}$$

We will prove this theorem in section 5. This will solve the control problem if we can check that $f(V)$ is indeed smaller than V. This will be proved in section 7. To see this, it is sufficient to note that in expression (2.1) of $F(W)$, the terms $\mathrm{e}^{-\{\Gamma W\}}$ and $(1-\mathrm{e}^{-\{\Gamma W\}})/\{\Gamma W\}$ are near $\mathbf{1}$ when $W \approx 0$. So

$$F(W) \approx \mathcal{R}W + \mathcal{N}W := W \qquad \text{so that} \quad f(W) = \mathcal{O}(W^2) \tag{2.5}$$

as is seen in (7.10).

Formula (2.3) connects the perturbed flow, modified by a control term, with the unperturbed flow.

The new factor, the flow of $\mathcal{R}V$ will turn out to commute with the flow of H, cf (5.9).

The second problem (the 'change of coordinates') is solved by an inversion formula. Let us rewrite (2.3) with $F(V)$ instead of $V + f(V)$, and with W instead of V

$$\forall t \in \mathbb{R} \qquad \mathrm{e}^{t\{H+F(W)\}} = \mathrm{e}^{-\{\Gamma W\}} \cdot \mathrm{e}^{t\{H\}} \cdot \mathrm{e}^{t\{\mathcal{R}W\}} \cdot \mathrm{e}^{\{\Gamma W\}}. \tag{2.6}$$

Theorem 2. *If we can find* W *such that* $F(W) = V$, i.e. $W = F^{-1}(V)$, *then*

$$\forall t \in \mathbb{R} \qquad \mathrm{e}^{t\{H+V\}} = \mathrm{e}^{-\{\Gamma W\}} \cdot \mathrm{e}^{t\{H\}} \cdot \mathrm{e}^{t\{\mathcal{R}W\}} \cdot \mathrm{e}^{\{\Gamma W\}}. \tag{2.7}$$

Here again, the flow of $\mathcal{R}W$ commutes with the flow of H.

Hence we need to invert the function F, but F is near the identity function around 0. More precisely we will find a ball around 0 in $\mathcal{A}$, for some norm, such that the difference between F and the identity function (what we called f in (2.2)) is Lipschitz and contractant, cf (2.5). So F can be inverted, at least around 0 (see section 7).

So to summarize this introduction, the first problem is solved explicitly for any perturbation, but we need to assume some smallness condition on the size of the perturbation to ensure that the 'control' term is smaller than the original perturbation. And the second problem is solved by an inversion formula, which also needs some smallness condition on the size of the perturbation to ensure that F is invertible. To be more precise, we need some norm on the Lie-algebra $\mathcal{A}$; this is done in section 7.

The equation in the unknown W (i.e. the inversion $W = F^{-1}(V)$) may be named the 'Hamilton–Jacobi equation'. Indeed, it yields the automorphism $\mathrm{e}^{\{\Gamma W\}}$.

Remark 2. If we replace hypothesis 1 by a (slightly) stronger one, then the situation is clearer, and simpler:

We assume that there exists a linear operator $\Gamma : \mathcal{A} \to \mathcal{A}$ such that

$$\{H\}^2\Gamma = \{H\}\Gamma\{H\} = \{H\} \tag{2.8}$$

$$\{H\}\Gamma^2 = \Gamma\{H\}\Gamma = \Gamma. \tag{2.9}$$

And we define four operators:

$$\mathcal{N} := \{H\}\Gamma \qquad \mathcal{R} := \mathbf{1} - \mathcal{N} \tag{2.10}$$

$$\tilde{\mathcal{N}} := \Gamma\{H\} \qquad \tilde{\mathcal{R}} := \mathbf{1} - \tilde{\mathcal{N}}. \tag{2.11}$$

Under this stronger assumption, we easily show that

$$\mathcal{N}^2 = \mathcal{N} \qquad \mathcal{R}^2 = \mathcal{R} \qquad \tilde{\mathcal{N}}^2 = \tilde{\mathcal{N}} \qquad \tilde{\mathcal{R}}^2 = \tilde{\mathcal{R}} \tag{2.12}$$

$$\operatorname{Ker}\tilde{\mathcal{R}} = \operatorname{Rg}\tilde{\mathcal{N}} = \operatorname{Rg}\Gamma \subset \operatorname{Ker}\mathcal{R} = \operatorname{Rg}\mathcal{N} = \operatorname{Rg}\{H\} \tag{2.13}$$

$$\operatorname{Ker}\mathcal{N} = \operatorname{Rg}\mathcal{R} = \operatorname{Ker}\Gamma \subset \operatorname{Ker}\tilde{\mathcal{N}} = \operatorname{Rg}\tilde{\mathcal{R}} = \operatorname{Ker}\{H\} \tag{2.14}$$

so that we have a characterization of the sub-Lie-algebra $\operatorname{Ker}\{H\}$: the 'constants of the motion' are exactly $\operatorname{Rg}\tilde{\mathcal{R}}$.

Furthermore, the above two 'set-inequalities' become 'set-equalities' if and only if

$$\{H\}\Gamma = \Gamma\{H\} \qquad \text{i.e. } \tilde{\mathcal{N}} = \mathcal{N}. \tag{2.15}$$

In that case (2.8) and (2.9) are equivalent to (1.6). For instance, assumption (2.15) is satisfied in examples (1.16) or (1.28).

But we do not need this stronger hypothesis for the rest of this paper.

3. Localization of the action variable

Let us study the case of classical mechanics, as above, with action-angle variables. Usually, we localize the action variable A near some point A_0 i.e. we make the canonical change of variables, from (A, θ) to (A_1, θ_1) with:

$$A = A_0 + \varepsilon A_1 \qquad \theta = \theta_1 \tag{3.1}$$

along with a rescaling of the Hamiltonian

$$H_1(A) := H(A_0 + \varepsilon \cdot A)/\varepsilon \tag{3.2}$$

where we choose some positive constant ε. This transformation is called a 'canonical similarity' since it preserves the symplectic form, up to the multiplicative constant ε. When H is integrable in the usual sense, with action-angle variables, we can expand the Hamiltonian $H = H(A)$ around A_0:

$$\begin{aligned} &H(A_0 + \varepsilon A_1) = c + \varepsilon \cdot \omega \cdot A_1 + q(\varepsilon \cdot A_1) \qquad \text{where} \\ &\omega := H'(A_0) \quad \text{and} \quad q(0) = q'(0) = 0 \end{aligned} \tag{3.3}$$

(i.e. q is quadratic in A_1) and where the additive constant $c := H(A_0)$ is not relevant and will be forgotten. Hence

$$H_1(A) = \omega \cdot A + \varepsilon \cdot \tilde{q}(A) \tag{3.4}$$

where $\tilde{q}(A) := q(\varepsilon \cdot A)/\varepsilon^2$ is of order ε^0.

Let us introduce a perturbation V, as above, and choose $\varepsilon := \|V\|^{1/2}$ for some norm. Under the rescaling (3.2) of the Hamiltonian, the perturbation is also divided by ε and so becomes

$$V_1(A, \theta) := V(A_0 + \varepsilon \cdot A, \theta)/\varepsilon \tag{3.5}$$

which is of order ε, so that the perturbed Hamiltonian is now

$$H_1(A) + V_1(A, \theta) = \omega \cdot A + \varepsilon \cdot V_2(A, \theta) \tag{3.6}$$

with

$$V_2(A, \theta) = \tilde{q}(A) + V_1(A, \theta)/\varepsilon \tag{3.7}$$

which is of order ε^0. So to summarize, it is always possible to assume that the integrable part is a harmonic oscillator, at least locally in the variable A, i.e. in a region (around any fixed A_0) where A (in the Hamiltonian (3.6)) is of order ε^0, but it may be wrong when A is of order $1/\varepsilon$. Hence in this case of classical mechanics, localized in action variable, the operator Γ is

$$\Gamma = \frac{1}{\omega \cdot \partial} \cdot \mathcal{N} \qquad \text{with} \quad \partial := \partial_\theta. \tag{3.8}$$

Let us express the action of Γ on an arbitrary trigonometric observable:

$$\Gamma \mathrm{e}^{\mathbf{i}\theta\Delta} = \frac{\mathrm{e}^{\mathbf{i}\theta\Delta}}{\mathbf{i}\omega \cdot \Delta} \cdot \chi(\Delta \in \mathbb{Z}^L \backslash \omega^\perp). \tag{3.9}$$

So that the operator $\mathcal{N}$ given by (1.15) after a Fourier transformation from the angles θ to the integer vector Δ, becomes

$$(\mathcal{N} V)(A, \Delta) = V(A, \Delta) \cdot \chi(\Delta \in \mathbb{Z}^L \backslash \omega^\perp). \tag{3.10}$$

4. Non-resonant Hamiltonians

We can also define the notion of 'non-resonance' as follows:

Definition 1. *In classical mechanics, H is 'non-resonant' iff:*

$$\forall\, V, W \in \operatorname{Ker}\{H\} \qquad \text{we have } \{V\}W = 0\cdot \tag{4.1}$$

In quantum mechanics, H is 'non-resonant' iff:

$$\forall\, V, P \in \operatorname{Ker}\{H\} \text{ s.t. } P^2 = P \qquad \text{we have } \{V\}P = 0. \tag{4.2}$$

Hence in the classical case

$$\forall\, W, V \in \mathcal{A} \quad \{\mathcal{R}W\}\mathcal{R}V = 0 \qquad \text{i.e. } \{\mathcal{R}\mathcal{A}\}\mathcal{R} = 0. \tag{4.3}$$

The above definition means that any two constants of motion of H do commute together.

After the localization of the action variable, as in section 3, the above non-resonant condition (4.1) is equivalent to the 'usual' non-resonance condition, which is in classical mechanics

$$\omega^\perp \cap \mathbb{Z}^L = \{0\} \tag{4.4}$$

i.e. there is no integer vector Δ orthogonal to ω, except for $\Delta = 0$. And the operator $\mathcal{R}$ may be written as the multiplication by the characteristic function $\chi(\Delta = 0)$.

We will give an example of a resonant Hamiltonian in (9.1), for which we can still apply our control theory.

In quantum mechanics, the non-resonance condition (4.2) means that any spectral projector of H commutes with any constant of motion. The non-resonance condition is 'usually' defined for the Floquet case

$$H = \sum_{k\in\mathbb{Z}, A\in\mathbb{N}} (k + h(A)) \cdot P_{k,A} \tag{4.5}$$

for some projectors $P_{k,A}$ which are mutually orthogonal. For instance the function $h(A)$ may be taken as $\omega \cdot A$ or $\omega \cdot A^2$. Condition (4.2) is satisfied exactly when the set of the spectral gaps intersects $\mathbb{Z}$ in the only point 0

$$\mathcal{G} \cap \mathbb{Z} = \{0\} \qquad \text{where} \quad \mathcal{G} := \{h(A) - h(B) \quad \text{s.t. } A, B \in \mathbb{N}\}. \tag{4.6}$$

For the case where $h(A) = \omega \cdot A^a$ for some positive integer a, condition (4.6) exactly means that the 'frequency' ω is 'non-resonant' in the usual sense. We see that the 'resonance' condition is a different notion than the degeneracy property, i.e. the dimension of the spectral projectors.

Let us note that the r.h.s. of (2.7) is called the 'normal form' of the (perturbed) flow on the l.h.s. When H is 'resonant', then (2.7) is called the 'resonant normal form'.

5. Proof of theorems 1 and 2

Actually theorems (2.3) and (2.7) are two different interpretations of the same formula (2.3) or (2.6). Let us first prove the following.

Proposition 1.

$$\forall\, W \in \mathcal{A} \qquad H + F(W) = \mathrm{e}^{-\{\Gamma W\}}(H + \mathcal{R}W). \tag{5.1}$$

Proof. Indeed from definitions (2.1) and (1.7), $F(W)$ can be rewritten as

$$F(W) - \mathrm{e}^{-\{\Gamma W\}}\mathcal{R}W = \frac{1 - \mathrm{e}^{-\{\Gamma W\}}}{\{\Gamma W\}}\{H\}\Gamma W = -\frac{1 - \mathrm{e}^{-\{\Gamma W\}}}{\{\Gamma W\}}\{\Gamma W\}H = \mathrm{e}^{-\{\Gamma W\}}H - H \tag{5.2}$$

where we used the antisymmetry (1.2). □

Proof of 2.6. Let us now take the brackets of the two sides of (5.1):

$$\{H + F(W)\} = \{\mathrm{e}^{-\{\Gamma W\}}(H + \mathcal{R}W)\}. \tag{5.3}$$

But

$$\forall\, V, W \in \mathcal{A} \qquad \{\mathrm{e}^{\{V\}}W\} = \mathrm{e}^{\{V\}} \cdot \{W\} \cdot \mathrm{e}^{-\{V\}}. \tag{5.4}$$

Indeed

$$\forall\, V, W \in \mathcal{A} \quad \forall\, n \in \mathbb{N} \qquad \{\{V\}^n W\} = \sum_{k=0}^{n} \binom{n}{k} \cdot \{V\}^{n-k} \cdot \{W\} \cdot \{-V\}^k. \tag{5.5}$$

The proof of (5.5) is an easy recurrence from the case $n = 1$, cf (1.3). Hence

$$\{H + F(W)\} = \mathrm{e}^{-\{\Gamma W\}} \cdot \{H + \mathcal{R}W\} \cdot \mathrm{e}^{\{\Gamma W\}}. \tag{5.6}$$

Let us now exponentiate the two sides of (5.6) (multiplied by any $t \in \mathbb{R}$):

$$\mathrm{e}^{t\{H+F(W)\}} = \exp[t \cdot \mathrm{e}^{-\{\Gamma W\}} \cdot \{H + \mathcal{R}W\} \cdot \mathrm{e}^{\{\Gamma W\}}] = \mathrm{e}^{-\{\Gamma W\}} \cdot \mathrm{e}^{t\{H+\mathcal{R}W\}} \cdot \mathrm{e}^{\{\Gamma W\}} \tag{5.7}$$

where we have used

$$\forall\, A, B \qquad \mathrm{e}^{A^{-1} \cdot B \cdot A} = A^{-1} \cdot \mathrm{e}^B \cdot A. \tag{5.8}$$

To conclude the proof of (2.6), it remains to show that $\{H\}$ and $\{\mathcal{R}W\}$ commute in the algebra of linear operators on $\mathcal{A}$. But from (1.3) and (1.10):

$$\{H\} \cdot \{\mathcal{R}W\} - \{\mathcal{R}W\} \cdot \{H\} = \{\{H\}\mathcal{R}W\} = 0. \tag{5.9}$$

Formula (2.6) is proved, and also theorems 1 and 2. □

6. The two main applications

6.1. Application 1: classical mechanics

Here, an automorphism is called a 'canonical transformation'.

First we have to express the Poisson bracket:

$$\forall\,(A,\Delta)\in\mathcal{H}\times\mathbb{Z}^L$$
$$\begin{aligned}(\{W\}V)(A,\Delta) = \mathbf{i}\sum_{\Delta'\in\mathbb{Z}^L} &(W'(A,\Delta-\Delta')\cdot\Delta')\cdot V(A,\Delta')\\ &-(V'(A,\Delta-\Delta')\cdot\Delta')\cdot W(A,\Delta')\end{aligned}\tag{6.1}$$

where W' is the derivative (gradient) of W with respect to A, and we have used a scalar product $W'(A,\Delta-\Delta')\cdot\Delta'$. If H is 'non-resonant', the action of the operator $\mathcal{R}$ on any observable W consists in keeping (in the Fourier coefficients of W) the only term with $\Delta=0$, i.e. the average over the angle variables θ. So $\mathcal{R}W$ is a function of A only (i.e. also 'integrable'), and any two functions of A only, do commute mutually. Then the flow $\mathrm{e}^{t\{\mathcal{R}W\}}$ is of the same type as the flow $\mathrm{e}^{t\{H\}}$, which is

$$\mathrm{e}^{t\{H\}} = \mathrm{e}^{t\cdot\omega(A)\partial_\theta}.\tag{6.2}$$

This is the operator translating the variable θ by $t\cdot\omega(A)$:

$$(\mathrm{e}^{t\{H\}}V)(A,\theta) = V(A,\theta+t\cdot\omega(A))\tag{6.3}$$

or after a Fourier transformation in θ:

$$(\mathrm{e}^{t\{H\}}V)(A,\Delta) = \mathrm{e}^{\mathbf{i}t\cdot\omega(A)\cdot\Delta}\cdot V(A,\Delta).\tag{6.4}$$

And similarly for $\mathrm{e}^{t\{\mathcal{R}W\}}$, with $\omega(A)$ replaced by $(\mathcal{R}W)'(A)=(\mathcal{R}W')(A)$.

6.2. Application 2: quantum mechanics

When $\mathcal{A}$ has an involution ('$*$', named 'adjonction', cf remark 1), and when $H=H^*$ the associated automorphism is called a 'unitary transformation'.

The bracket is defined as

$$\forall\,A\in\mathcal{H}\qquad\forall\Delta\in\mathcal{G}(A)$$
$$\begin{aligned}(\{W\}V)(A,\Delta) = \frac{\mathbf{i}}{\hbar}\sum_{\Delta'\in\mathcal{G}(A)} &W(A+\Delta',\Delta-\Delta')\cdot V(A,\Delta')\\ &-V(A+\Delta',\Delta-\Delta')\cdot W(A,\Delta').\end{aligned}\tag{6.5}$$

We can put this bracket in a form similar to the Poisson bracket (6.1), by adding and subtracting two terms:

$$\forall\,A\in\mathcal{H}\qquad\forall\,\Delta\in\mathcal{G}(A)$$
$$\begin{aligned}(\{W\}V)(A,\Delta) = \sum_{\Delta'\in\mathcal{G}(A)} &\mathbf{i}\left[\frac{W(A+\Delta',\Delta-\Delta')-W(A,\Delta-\Delta')}{\hbar}\right]\cdot V(A,\Delta')\\ &-\mathbf{i}\left[\frac{V(A+\Delta',\Delta-\Delta')-V(A,\Delta-\Delta')}{\hbar}\right]\cdot W(A,\Delta')\\ &+\{W(A,\Delta')\}V(A,\Delta-\Delta')\end{aligned}\tag{6.6}$$

where the last term is a short notation for

$$\mathbf{i}\left[\frac{W(A,\Delta')\cdot V(A,\Delta-\Delta')-V(A,\Delta-\Delta')\cdot W(A,\Delta')}{\hbar}\right].\tag{6.7}$$

Formula (6.6) is reminiscent of the famous ‘correspondence principle’ between classical mechanics and quantum mechanics: if the set $\mathcal{H}$ becomes more and more dense in $\mathbb{R}$ (for instance if it is $\hbar\mathbb{Z}$) with $\hbar \to 0$, in such a way that A remains constant, then the last term (6.7) tends to 0 and the first term in the bracket tends to the derivative of W with respect to A, multiplied by Δ', so that the quantum bracket becomes the classical one; cf [1] for a precise formulation of this fact, in some particular cases.

From now on, we will choose $\hbar = 1$, since we will not use this semi-classical limit.

Here again, if H is ‘non-resonant’, $\mathcal{R}W$ is a diagonal matrix, and so its flow is of the same type as the flow of H, since:

$$(\mathrm{e}^{t\{H\}}V)(A,\Delta) = \mathrm{e}^{\mathbf{i}t\cdot[h(A+\Delta)-h(A)]} \cdot V(A,\Delta) \tag{6.8}$$

and similarly for the other flow.

7. Quantitative estimates

We start by choosing an arbitrary norm on $\mathcal{A}$, and we replace $\mathcal{A}$ by its closure with respect to this norm. We deduce a canonical norm for the operator Γ which acts bilinearly on $\mathcal{A}$:

$$|||\Gamma||| := \sup_{V,W\in\mathcal{A}\ \mathrm{s.t.}\ \|V\|=\|W\|=1} \|\{\Gamma W\}V\|. \tag{7.1}$$

We make an important assumption.

Hypothesis 2.

$$|||\Gamma||| < \infty. \tag{7.2}$$

This hypothesis is necessary to apply the so-called ‘local bijection theorem’ to invert the function F. When (7.2) is not true, it can be replaced by a weaker one, but then we need to use the more complicated theorem of Nash–Moser, cf [7, 9, 11, 12] or the Newton iterative method, as in the KAM theory. They are based on a Frechet structure on $\mathcal{A}$ i.e. an infinite sequence of norms $\|\cdots\|_{s-1} \geqslant \|\cdots\|_s$ instead of only 1 norm. Hypothesis (7.2) is still required but with a weaker norm:

$$|||\Gamma|||_\alpha := \sup_{s\in\mathbb{N}} \sup_{V,W\in\mathcal{A}\ \mathrm{s.t.}\ \|V\|_{s-1}=\|W\|_s=1} \|\{\Gamma V\}W\|_s \cdot \alpha(s) \tag{7.3}$$

where $\alpha : \mathbb{N} \to \mathbb{R}_+^*$ and $\lim_{s\to\infty} \alpha(s) = 0$. Hence Γ may be bounded if we admit some loss of ‘regularity’.

Let us also define a norm

$$||||\mathcal{R}|||| := \sup_{W\in\mathcal{A}\ \mathrm{s.t.}\ \|W\|=1} \|\mathcal{R}W\|. \tag{7.4}$$

Another assumption is on the operator $\mathcal{R}$.

Hypothesis 3.

$$||||\mathcal{R}|||| \leqslant 1. \tag{7.5}$$

It is fulfilled for many norms, for instance those given in (10.1), cf (10.2).

Then F is invertible when $\|V\|$ is small enough.

Theorem 3. *Let V be an element of $\mathcal{A}$ and Γ be defined by (1.6), or explicitly by (1.16) or (1.28). Under the hypotheses 1–3:*

$$\textit{if}\ \|V\| \leqslant \frac{1}{5|||\Gamma|||} \tag{7.6}$$

$$then\ \frac{24}{35} \leqslant \frac{\|F^{-1}(V)\|}{\|V\|} \leqslant \frac{24}{13}. \tag{7.7}$$

Proof. Let us start by expanding $f(W)$ as given by (2.2) in series

$$f(W) = \sum_{n=1}^{\infty} \{-\Gamma W\}^n \cdot \frac{n\mathcal{R}+1}{n+1!} \cdot W \tag{7.8}$$

and use definition (7.1):

$$\|\{\Gamma W\}V\| \leqslant |||\Gamma||| \cdot \|W\| \cdot \|V\| \tag{7.9}$$

and hypothesis (7.5), so that

$$\|f(W)\| \leqslant \sum_{n=1}^{\infty} (|||\Gamma||| \cdot \|W\|)^n \cdot \frac{n||||\mathcal{R}|||| + 1}{n+1!} \cdot \|W\| \leqslant (\mathrm{e}^{|||\Gamma|||\cdot\|W\|} - 1) \cdot \|W\|. \tag{7.10}$$

This proves that $f(W) = \mathcal{O}(W^2)$. To compute the derivative of $f(W)$ with respect to W we need the following formula, valid for any derivation $\partial \in \mathrm{Der}(\mathcal{L})$ of some algebra $\mathcal{L}$:

$$\partial\,\mathrm{e}^V = \int_0^1 \mathrm{d}t \cdot \mathrm{e}^{t\cdot V} \cdot \partial V \cdot \mathrm{e}^{(1-t)\cdot V} \tag{7.11}$$

which may be rewritten as

$$\partial\,\mathrm{e}^V = \left(\frac{\mathrm{e}^{\{V\}} - 1}{\{V\}} \partial V\right) \cdot \mathrm{e}^V \tag{7.12}$$

where the Lie-bracket $\{\cdots\}$ is given by the commutator. Indeed

$$\mathrm{e}^V \cdot W \cdot \mathrm{e}^{-V} = \mathrm{e}^{\{V\}} W. \tag{7.13}$$

The proof of (7.12) is obtained by power expanding the exponentials

$$\forall n \in \mathbb{N} \qquad \partial(V^n) = \sum_{k=1}^{n} \binom{n}{k} \cdot \{V\}^{k-1} \cdot (\partial V) \cdot V^{n-k} \tag{7.14}$$

which is proved by a simple recurrence. Similarly for (7.11)

$$\partial\,\mathrm{e}^V = \sum_{N=0}^{\infty} \frac{\partial(V^N)}{N!} = \sum_{n,k=0}^{\infty} \frac{V^n \cdot \partial V \cdot V^k}{(n+k+1)!} \tag{7.15}$$

whereas

$$\int_0^1 \mathrm{d}t \cdot \mathrm{e}^{t\cdot V} \cdot \partial V \cdot \mathrm{e}^{(1-t)\cdot V} = \sum_{n,k=0}^{\infty} \int_0^1 \mathrm{d}t \cdot \frac{t^n \cdot (1-t)^k}{n! \cdot k!} V^n \cdot \partial V \cdot V^k. \tag{7.16}$$

These two expressions coincide after using:

$$\int_0^1 \mathrm{d}t \cdot \frac{t^n \cdot (1-t)^k}{n! \cdot k!} = \frac{1}{(n+k+1)!}. \tag{7.17}$$

Let us note that (7.11) is a generalization of the following formula (valid when t varies in a finite set) to the case where t is a continuous variable:

$$\partial\left(\prod_t V_t\right) = \sum_t \left(\prod_{\tau<t} V_\tau\right)(\partial V_t)\left(\prod_{\tau>t} V_\tau\right). \tag{7.18}$$

When ∂V commute with V we retrieve the usual formula:

$$\partial\,\mathrm{e}^V = (\partial V) \cdot \mathrm{e}^V = \mathrm{e}^V \cdot (\partial V). \tag{7.19}$$

Let us apply (7.11) to the Lie algebra $\mathcal{L} = \mathcal{L}(\mathcal{A})$ of the endomorphisms of $\mathcal{A}$ (the space of observables), for which the bracket is indeed the commutator. And we take for ∂ the derivation with respect to W:

$$\partial\, \mathrm{e}^{\{\Gamma W\}} = \int_0^1 \mathrm{d}t \cdot \mathrm{e}^{t\cdot\{\Gamma W\}} \cdot \{\Gamma \ldots\} \cdot \mathrm{e}^{(1-t)\cdot\{\Gamma W\}} \tag{7.20}$$

which can be rewritten as

$$\partial\, \mathrm{e}^{\{\Gamma W\}} = \left(\frac{\mathrm{e}^{\{\{\Gamma W\}\}} - 1}{\{\{\Gamma W\}\}} \cdot \{\Gamma \ldots\}\right) \cdot \mathrm{e}^{\{\Gamma W\}} \tag{7.21}$$

where the double bracket is the bracket in $\mathcal{L}(\mathcal{A})$ i.e. the commutator:

$$\forall\, X \in \mathcal{L}(\mathcal{A}) \qquad \{\{\Gamma W\}\}X := \{\Gamma W\} \cdot X - X \cdot \{\Gamma W\}. \tag{7.22}$$

Therefore

$$\begin{aligned} F'(W)V &= (\psi(\{\{\Gamma W\}\})\{\Gamma V\})\,\mathrm{e}^{-\{\Gamma W\}}\mathcal{R}W + \mathrm{e}^{-\{\Gamma W\}}\mathcal{R}V \\ &\quad + \left(\frac{1 - \mathrm{e}^{-\{\Gamma W\}}}{\{\Gamma W\}}\right)\mathcal{N}V + (\varphi(\{\{\Gamma W\}\})\{\Gamma V\})\,\mathrm{e}^{-\{\Gamma W\}}\mathcal{N}W \end{aligned} \tag{7.23}$$

where

$$\psi(X) := \frac{1 - \mathrm{e}^{-X}}{X} \qquad \varphi(X) := \int_0^1 \mathrm{d}s \cdot \psi(sX). \tag{7.24}$$

Let us now use formula (7.9):

$$\|\mathrm{e}^{\{\Gamma W\}} \cdot V\| \leqslant \mathrm{e}^{|||\Gamma|||\cdot\|W\|} \cdot \|V\| \tag{7.25}$$

so that

$$\|\partial\, \mathrm{e}^{\{\Gamma W\}} \cdot V\| \leqslant \int_0^1 \mathrm{d}t \cdot \mathrm{e}^{t\cdot|||\Gamma|||\cdot\|W\|} \cdot |||\Gamma||| \cdot \mathrm{e}^{(1-t)\cdot|||\Gamma|||\cdot\|W\|} \cdot \|V\| = \mathrm{e}^{|||\Gamma|||\cdot\|W\|} \cdot |||\Gamma|||.\|V\|. \tag{7.26}$$

And we do not need formula (7.23). Hence $f'(W)$ has a norm bounded by

$$\begin{aligned} \|f'(W)\| \leqslant \sum_{n=1}^{\infty} & n \cdot (|||\Gamma||| \cdot \|W\|)^{n-1} \cdot |||\Gamma||| \cdot \frac{n||||\mathcal{R}|||| + 1}{n+1!} \cdot \|W\| \\ & + (|||\Gamma||| \cdot \|W\|)^n \cdot \frac{n||||\mathcal{R}|||| + 1}{n+1!} \end{aligned} \tag{7.27}$$

so that, using (7.5):

$$\|f'(W)\| \leqslant \mathrm{e}^{|||\Gamma|||\cdot\|W\|} \cdot (|||\Gamma|||\cdot\|W\| + 1) - 1. \tag{7.28}$$

Let us call γ the solution of the transcendental equation:

$$\mathrm{e}^{\gamma}(\gamma + 1) - 1 = 1 \qquad \text{i.e. } \gamma = 0.374\,822\,525\,811\,8948\cdots > 13/35. \tag{7.29}$$

Hence

$$\text{if} \quad |||\Gamma|||\cdot\|W\| < \gamma \qquad \text{then} \quad \|f'(W)\| < 1 \tag{7.30}$$

and so

$$\|F'(W)\| = \|1 + f'(W)\| \geqslant 1 - \|f'(W)\| > 0. \tag{7.31}$$

Then F is invertible and

$$\|F(W)\| = \|W + f(W)\| \leqslant \|W\| \cdot \left(1 + \frac{\|f(W)\|}{\|W\|}\right) \leqslant \|W\| \cdot\ (1 + \mathrm{e}^{\gamma} - 1) = \mathrm{e}^{\gamma} \cdot \|W\|. \tag{7.32}$$

Likewise

$$\|F(W)\| \geqslant \|W\| \cdot \left(1 - \frac{\|f(W)\|}{\|W\|}\right) \geqslant \|W\| \cdot (2 - \mathrm{e}^{\gamma}). \tag{7.33}$$

So if we replace W by $F^{-1}(V)$ we get

$$\mathrm{e}^{-\gamma} \leqslant \frac{\|F^{-1}(V)\|}{\|V\|} \leqslant \frac{1}{2 - \mathrm{e}^{\gamma}} \tag{7.34}$$

under condition (7.30), i.e. if

$$\|V\| < \frac{(2 - \mathrm{e}^{\gamma}) \cdot \gamma}{|||\Gamma|||}. \tag{7.35}$$

Indeed we will have in that case:

$$\|W\| = \|F^{-1}(V)\| \leqslant \frac{\|V\|}{2 - \mathrm{e}^{\gamma}} < \frac{\gamma}{|||\Gamma|||}. \tag{7.36}$$

To prove theorem 3, it remains to use the value of γ: $\mathrm{e}^{\gamma} = \frac{2}{\gamma+1} < \frac{35}{24}$. □

Let us note that under condition (7.6), the new term (the 'control term', $f(V)$) will be smaller than V (cf (7.10)):

$$\text{if} \quad \|V\| \leqslant \frac{1}{5|||\Gamma|||} \qquad \text{then} \quad \frac{\|f(V)\|}{\|V\|} \leqslant \mathrm{e}^{1/5} - 1 \left(< \frac{2}{9}\right). \tag{7.37}$$

8. A formal series for the inverse of *F*

We want an expansion of $W = F^{-1}(V)$ in powers of V, using the expansion of $F(W) = W + f(W)$ given in (7.8). For that purpose we can use an extension of the Lagrange inversion formula which was established to invert a function from $\mathbb{C} \to \mathbb{C}$ in power series. An extension to the case where the argument is not a number but a function (as given here), was given in [14], and later in [2]. We first rewrite the definition of W as a fixed-point problem:

$$F(W) = V \quad \Longleftrightarrow \quad W = G(W) := V - f(W). \tag{8.1}$$

A Taylor expansion of G is given by

$$G(W) = \sum_{n \geqslant 0} \hat{G}(n) W^n \tag{8.2}$$

with (cf 7.8)

$$\hat{G}(0) := G(0) = V \qquad \text{and} \qquad \hat{G}(1) := 0 \tag{8.3}$$

$$\forall n \geqslant 2 \qquad \hat{G}(n) := \frac{G^{(n)}(0)}{n!} = \mathcal{S}(-1)^n \{\Gamma \ldots\}^{n-1} \cdot \frac{(n-1)\mathcal{R} + 1}{n!} \ldots \tag{8.4}$$

is an n-linear completely symmetric application from $\mathcal{A}^n$ to $\mathcal{A}$. Its n arguments are symbolized by the n 'slots'. In (8.2), this n-linear mapping is applied to identical arguments, n times W. When applied to n general arguments $W_1, \ldots, W_n$, it would give

$$\hat{G}(n)(W_1, \ldots, W_n) = \mathcal{S}(-1)^n \{\Gamma W_1\} \ldots \{\Gamma W_{n-1}\} \cdot \frac{(n-1)\mathcal{R} + 1}{n!} W_n. \tag{8.5}$$

The operator $\mathcal{S}$ is the symmetrization of the arguments which yields, when applied to an n-linear application T,

$$(\mathcal{S}T)(W_1, \ldots, W_n) := \frac{1}{n!} \sum_{\sigma \in \text{Permutations}} T(W_{\sigma(1)}, \ldots, W_{\sigma(n)}). \tag{8.6}$$

We will omit the parentheses and the commas in using such tensors. So $W^M = (W, W, \dots, W)$ (M times). For instance

$$G'(0)(W_1) := \lim_{\lambda\to 0} \frac{G(\lambda W_1) - G(0)}{\lambda} \tag{8.7}$$

which is 0, for the function defined in (8.1), since it is quadratic in its argument. And $G''(0)$ is a tensor of order 2:

$$G''(0)(W_1, W_2) := \lim_{\lambda\to 0} \lim_{\mu\to 0} \frac{G(\lambda W_1 + \mu W_2) - G(\lambda W_1) - G(\mu W_2) + G(0)}{\lambda \cdot \mu} \tag{8.8}$$

which can be easily computed to be

$$\hat{G}(2) W_1 \cdot W_2 = \{\Gamma W_1\} \frac{\mathcal{R}+1}{4} W_2 + \{\Gamma W_2\} \frac{\mathcal{R}+1}{4} W_1. \tag{8.9}$$

Theorem 4. *The above-mentioned extension of the Lagrange inversion formula says that the solution of any fixed-point equation $W = G(W)$, with a function G given by a series (8.2), is formally:*

$$W = \sum_{N\geqslant 1} \sum_{\nu\in\mathcal{T}(N)} \hat{G}(\nu_{N-1})\hat{G}(\nu_{N-2}) \cdots \hat{G}(\nu_1)\hat{G}(\nu_0) \tag{8.10}$$

where

$$\mathcal{T}(N) := \{\nu = (\nu_0, \nu_1, \dots, \nu_{N-1}) \in \{0, 1, \dots, N-1\}^N \text{s.t.}$$
$$\forall k \in \{0, 1, \dots, N-1\} : |\nu|_k \leqslant k \qquad \text{and} \qquad |\nu|_{N-1} = N-1\} \tag{8.11}$$

with

$$|\nu|_k := \nu_0 + \nu_1 + \cdots + \nu_k. \tag{8.12}$$

The expansion (8.10) is only useful if $\hat{G}(n)$ are 'small', since we expand in their powers. In our cases (8.3) and (8.4), only $\hat{G}(0)$ is small, and $\hat{G}(1) = 0$. But $\hat{G}(n)$ is of order 1, when $n \geqslant 2$. So we have to rearrange the series (8.10):

$$W = V + \sum_{M\geqslant 2} W_M \tag{8.13}$$

where

$$W_M := \sum_{N=M+1}^{2M-1} \sum_{\nu\in\mathcal{T}(N)\,\text{s.t.}\,\nu^{-1}(0)=M} \hat{G}(\nu_{N-1})\hat{G}(\nu_{N-2}) \cdots \hat{G}(\nu_1)\hat{G}(\nu_0) \tag{8.14}$$

with the restriction '$\nu^{-1}(0) = M$' in the sum over ν, meaning that ν must take exactly M times the value 0. Indeed this is exactly the case '$\nu_k = 0$' which produces a factor V. We can also note that in definition (8.11) of $\mathcal{T}(N)$ we can restrict $\nu \in \{0, 2, 3, \dots, N-1\}$ since in our case $\hat{G}(1) = 0$. Before we explicit the first terms of the solution (8.13), let us understand heuristically the series (8.10) or (8.13). A first approximation of the solution of $W = G(W)$ is $W \approx G(0) := V$. Then a second approximation is $W \approx G(G(0)) \approx G(0) + G'(0) \cdot G(0)$. And a third one, at the order V^3, is $W \approx G(G(G(0))) \approx G[G(0) + G'(0) \cdot G(0)]$ i.e.

$$W \approx G(0) + G'(0) \cdot G(0) + G'(0) \cdot G'(0) \cdot G(0) + \tfrac{1}{2} \cdot G''(0) \cdot G(0) \cdot G(0) \tag{8.15}$$

and so on. We are building the series (8.10). Now we have to check that the general term of (8.10) is indeed an element of $\mathcal{A}$ and not an arbitrary m-linear mapping for some m. Each term is a composition of high-order tensors, and this product is actually a 'vector' i.e. an element of $\mathcal{A}$. This is due to the definition of $\mathcal{T}(N)$: we note that $\hat{G}(n)$ is of type $V.\bar{V}^n$ i.e.

it is a (sum of) tensorial product(s) of a vector V and of n covectors $\bar{V}$. More generally we say that a tensor is of type $V^m \cdot \bar{V}^n$ when it is n times covariant and m times contravariant. This is a convenient way of considering such tensors, in order to keep track of all subsequent contractions. Of course, the ordering is crucial: $V \cdot \bar{V}$ is a matrix (a tensor of order 2) whereas $\bar{V} \cdot V$ is a scalar (a scalar product). Each time a covector follows (on the left) a vector, we make the contraction, i.e. we reduce the tensorial product to a scalar. Actually we just want to give a meaning to the composition of multilinear applications. To this purpose we can define

$$\forall B \in \mathcal{L}(\mathcal{A}^n, \mathcal{A}) \qquad \forall C \in \mathcal{L}(\mathcal{A}^m, \mathcal{A}) \qquad B \cdot C \in \mathcal{L}(\mathcal{A}^{m+n-1}, \mathcal{A}) \tag{8.16}$$

by

$$(B \cdot C)(V_1, \ldots, V_m, V_{m+1}, \ldots, V_{m+n-1}) := B(C(V_1, \ldots, V_m), V_{m+1}, \ldots, V_{m+n-1}) \in \mathcal{A}. \tag{8.17}$$

Now we note that for any $\nu \in \mathcal{T}(N)$: $\nu_0 = 0$ since $|\nu|_0 = 0$. So the right-most factor in (8.10) is a vector: $\hat{G}(0) = V \in \mathcal{A}$.

Then the preceding one is $\hat{G}(\nu_1)$ where $\nu_1 \leqslant 1$ since $|\nu|_1 = \nu_0 + \nu_1$ has to be less than 1. Hence $\hat{G}(\nu_1)$ is either a vector (if $\nu_1 = 0$) or a matrix (if $\nu_1 = 1$). In the latter case the product $\hat{G}(1)\hat{G}(0)$ is a vector, and in the former case we have a 'bi-vector' which will be made into a vector by the higher-order terms: indeed if $\nu_1 = 0$ then there exists a $k \geqslant 2$ such that $\nu_k \geqslant 2$ since $|\nu|_{N-1}$ has to be $N - 1$. More precisely we note that the product of a tensor $V^a \cdot \bar{V}^b$ by a tensor $V^c \cdot \bar{V}^d$ is a tensor of type

$$V^a \cdot \bar{V}^b \cdot V^c \cdot \bar{V}^d = V^{a+c-\min(b,c)} \cdot \bar{V}^{b+d-\min(b,c)}. \tag{8.18}$$

So the product $\hat{G}(\nu_k) \cdots \hat{G}(\nu_0)$ is a tensor of type

$$\hat{G}(\nu_k) \cdots \hat{G}(\nu_0) = V^{\max(k-|\nu|_k, 1)} \cdot \bar{V}^0 \tag{8.19}$$

i.e. without any covector, since we assume $|\nu|_k \leqslant k \, \forall k$. Finally for $k = N - 1$, $|\nu|_{N-1}$ has to be $N - 1$, so (8.10) is indeed a sum of vectors.

The idea of the proof of (8.10) is to plug the r.h.s. of (8.10) into the expression (8.2) of $G(W)$ and to recognise that we get the same expansion as (8.10), i.e. this is also W.

A simple illustration (which is not of our type) of theorem 4 is when

$$G(W) := V \cdot \mathrm{e}^{\bar{Y} \cdot W} \tag{8.20}$$

for some $V \in \mathcal{A}$ and $\bar{Y} \in \mathcal{A}^* := \mathcal{L}(\mathcal{A}, \mathbb{R})$. In that simple case, theorem 4 gives the solution of the 'transcendental' equation in W:

$$W = V \cdot \mathrm{e}^{\bar{Y} \cdot W} \tag{8.21}$$

as

$$W = V \cdot \lambda(\bar{Y} \cdot V) \qquad \text{where} \quad \lambda(x) := \sum_{n \geqslant 0} \frac{(n+1)^{n-1}}{n!} x^n. \tag{8.22}$$

So that $\lambda(x)$ converges if $|x| \leqslant \mathrm{e}^{-1}$ (but may be extended if $x < -\mathrm{e}^{-1}$).

We can rearrange the series (8.14) as follows:

$$W_M := \sum_{N=1}^{M-1} \sum_{\lambda \in \mathcal{B}(N,M)} \sum_{\mu \in \mathcal{C}(N,\lambda)} \hat{G}\left(\lambda_1^{\mu}\right) \cdots \hat{G}\left(\lambda_{N+M}^{\mu}\right) \tag{8.23}$$

where $\hat{G}$ is defined in (8.3) and (8.4), and

$$\mathcal{B}(N, M) := \{\lambda \in \mathbb{N}_*^N \text{s.t.} |\lambda|_N = M - 1\} \tag{8.24}$$

and

$$\mathcal{C}(N, \lambda) := \{\mu \in \mathbb{N}_*^N \text{ s.t. } \forall\, 1 \leqslant n \leqslant N\; |\mu|_n \leqslant n + |\lambda|_{n-1}\} \tag{8.25}$$

with

$$|\mu|_n := \mu_1 + \cdots + \mu_n \tag{8.26}$$

and similarly for $|\lambda|_n$ with $|\lambda|_0 := 0$. Finally in (8.23) the integer λ_l^μ is defined for any $1 \leqslant l \leqslant N + M$ by

$$\begin{aligned} \lambda_l^\mu &= \lambda_n + 1 \qquad \text{if} \quad l = |\mu|_n \\ &= 0 \qquad\qquad \text{otherwise.} \end{aligned} \tag{8.27}$$

This means that $\lambda_l^\mu = 0$ when l is not an element of the set of the values $|\mu|_n$, when n varies in $\{1, \ldots, N\}$. This case occurs M times when l varies in $\{1, \ldots, N + M\}$. Otherwise when $l = |\mu|_n$ for some n, then $\lambda_l^\mu = \lambda_n + 1$, for this n.

Finally we can also rearrange the Lagrange series in our specific cases (8.3) and (8.4) into

$$W_M := \sum_{N=1}^{M-1} \sum_{\lambda\in\mathcal{B}(N,M)} \sum_{\nu\in\mathcal{D}(N,\lambda)} (-1)^{N+M+1} \{\Gamma \ldots\}^{\lambda_N} \frac{\lambda_N \mathcal{R} + 1}{\lambda_N + 1!} V^{\nu_N} \ldots \{\Gamma \ldots\}^{\lambda_1} \frac{\lambda_1 \mathcal{R} + 1}{\lambda_1 + 1!} V^{\nu_1} \tag{8.28}$$

where

$$\mathcal{D}(N, \lambda) := \{\nu \in \mathbb{N}^N \qquad \text{s.t. } |\nu|_N = M \qquad \text{and} \qquad \forall\, 1 \leqslant n \leqslant N \ |\nu|_n > |\lambda|_n\}. \tag{8.29}$$

Let us note that

$$\|\hat{G}(n)\| \leqslant \frac{|||\Gamma|||^{n-1}}{(n-1)!} \tag{8.30}$$

and so

$$\|W_M\| \leqslant c_M \cdot |||\Gamma|||^{M-1} \cdot \|V\|^M \tag{8.31}$$

for some positive constant c_M. Indeed: $\lambda_1 + \cdots + \lambda_N = M - 1$.

Let us explicit the first orders of the expansion of the solution W of our problems (8.1)–(8.4), as given by (8.13) or (8.23):

$$F^{-1}(V) = W = \sum_{M\geqslant 1} W_M \qquad \text{where} \quad W_1 = V \quad W_2 = \{\Gamma V\}\frac{\mathcal{R}+1}{2} V$$

$$W_3 = \left(\{\Gamma V\}\frac{\mathcal{R}+1}{2}\right)^2 V - \{\Gamma V\}^2 \frac{2\mathcal{R}+1}{6} V + \left\{\Gamma\left(\{\Gamma V\}\frac{\mathcal{R}+1}{2} V\right)\right\}\frac{\mathcal{R}+1}{2} V$$

$$\begin{aligned} W_4 = {} & \left(\{\Gamma V\}\frac{\mathcal{R}+1}{2}\right)^3 V + \{\Gamma V\}^3 \frac{3\mathcal{R}+1}{24} V - \{\Gamma V\}^2 \frac{2\mathcal{R}+1}{6} \{\Gamma V\}\frac{\mathcal{R}+1}{2} V \\ & - \left\{\Gamma\left(\{\Gamma V\}\frac{\mathcal{R}+1}{2} V\right)\right\}\frac{\mathcal{R}+1}{2}\{\Gamma V\}\frac{\mathcal{R}+1}{2} V \\ & - \left\{\Gamma\left(\{\Gamma V\}^2\frac{2\mathcal{R}+1}{6} V\right)\right\}\frac{\mathcal{R}+1}{2} V - \{\Gamma V\}\frac{\mathcal{R}+1}{2}\left\{\{\Gamma V\}^2\frac{2\mathcal{R}+1}{6} V\right\} \\ & + \{\Gamma V\}\frac{\mathcal{R}+1}{2}\left\{\Gamma\left(\{\Gamma V\}\frac{\mathcal{R}+1}{2} V\right)\right\}\frac{\mathcal{R}+1}{2} V \\ & + \left\{\Gamma\left(\left\{\Gamma\left(\{\Gamma V\}\frac{\mathcal{R}+1}{2} V\right)\right\}\frac{\mathcal{R}+1}{2} V\right)\right\}\frac{\mathcal{R}+1}{2} V \\ & + \left\{\Gamma\left(\left(\{\Gamma V\}\frac{\mathcal{R}+1}{2}\right)^2 V\right)\right\}\frac{\mathcal{R}+1}{2} V \\ & - \left\{\{\Gamma V\}\Gamma\left(\{\Gamma V\}\frac{\mathcal{R}+1}{2} V\right)\right\}\frac{2\mathcal{R}+1}{6} V \end{aligned} \tag{8.32}$$

and so on for $W_5 \ldots$. We have proved in theorem 3 that this series converges at least when

$$\|V\| \leqslant \frac{1}{5|||\Gamma|||}. \tag{8.33}$$

The paper [8] proves that a series similar to this one (the Lindstedt series, which is local in the variable A) converges. The key ingredient is the compensation between the terms of different signs.

Remark. *As we noted at the end of section 1, the solution W is independent of constant $\hbar$ since the operator Γ always appears with a bracket around it. And the pair $\{\Gamma \ldots\}$ is independent of $\hbar$.*

9. Example 1: classical control theory

We will study a simple model, introduced in [13] and described in [4, 5], of a charged particle in a plasma, in a Tokamak, which is a reactor for controlled thermonuclear fusion. We consider a section of the Tokamak as the phase space of a dynamical system with one degree of freedom, so that the particle has one degree of freedom (p, q), but it is embedded in a complicated electric field, depending on time. The fast motion given by the strong magnetic field has been averaged out. The Hamiltonian is

$$H(p, q, E, \tau) + V(p, q, E, \tau) \qquad \text{where} \quad H(p, q, E, \tau) = E \tag{9.1}$$

and

$$V(p, q, E, \tau) = \sum_{n,m,k\neq 0} \varepsilon^{(1)}_{n,m,k} \cdot \sin(nq + mp + k\tau). \tag{9.2}$$

The extended canonical coordinates are (E, τ) so that the motion of the new dynamical variable τ is trivial:

$$\mathrm{e}^{t\cdot\{H+V\}}\tau = \tau + t. \tag{9.3}$$

Here we have applied the flow to the observable $\tau : (p, q, E, \tau) \mapsto \tau$, and of course V is independent of the variable E, which is the variable canonically conjugate to τ. Let us note that in (9.2), the variable k must be integer (or at least away from 0) but the variables m, n may be integers or real numbers: in that case the sum over them should be replaced by an integral.

The Hamiltonian H is resonant since $\{H\} = \partial_\tau$ so that

$$(\mathcal{R}V)(p, q, \tau) = \oint \mathrm{d}\tau \cdot V(p, q, \tau) \tag{9.4}$$

is independent of τ. Indeed a Fourier transformation in τ yields

$$(\widehat{\mathcal{R}V})(p, q, E, k) = \hat{V}(p, q, E, 0) \cdot \chi(k = 0). \tag{9.5}$$

For the perturbation (9.2) we have $\mathcal{R}V = 0$. The action of the operator Γ is defined by

$$(\Gamma V)(p, q, \tau) = -\sum_{n,m,k\neq 0} \frac{\varepsilon^{(1)}_{n,m,k}}{k} \cdot \cos(nq + mp + k\tau). \tag{9.6}$$

Then the 'control term' $f(V)$ can be explicitly computed, cf (7.8) with $\mathcal{R}V = 0$:

$$f(V) = \sum_{s\geqslant 2} f_s \qquad \text{where} \quad f_s := \frac{\{-\Gamma V\}^{s-1}}{s!} V. \tag{9.7}$$

Hence

$$f(V)(p,q,\tau) = \sum_{n,m,k\neq 0} \varepsilon_{n,m,k} \cdot \sin(nq + mp + k\tau) \tag{9.8}$$

where

$$\varepsilon_{n,m,k} := \sum_{s\geqslant 2} (-1)^{s-1} \cdot \frac{\varepsilon^{(s)}_{n,m,k}}{s!} \tag{9.9}$$

with

$$\varepsilon^{(s)}_{n,m,k} := \sum_{N,M,K\neq 0} \frac{\varepsilon^{(1)}_{N,M,K}}{K} \cdot \left(\varepsilon^{(s-1)}_{N+n,M+m,K+k} - \varepsilon^{(s-1)}_{N-n,M-m,K-k}\right) \cdot (M\cdot n - N\cdot m). \tag{9.10}$$

The proof of (9.10) is based on

$$\begin{aligned}\{\cos(Nq + Mp + K\tau)\}\sin(nq + mp + k\tau)\\ = \tfrac{1}{2}\cdot(N.m - M.n)\cdot[\sin[(N+n)q + (M+m)p + (K+k)\tau]\\ + \sin[(N-n)q + (M-m)p + (K-k)\tau]]\end{aligned} \tag{9.11}$$

so that we can iterate and compute $\{\Gamma V\}^s V$.

A simple case where the control term can be computed explicitly is as follows. Let us choose

$$\varepsilon \in \mathbb{R}^* \qquad b > 1/\sqrt{2} \qquad \sigma = \pm 1 \qquad m, n \in \mathbb{R} \quad \text{s.t.}\, m \neq n. \tag{9.12}$$

And we take V as a particular case of (9.2), a sum of two waves:

$$V(p,q,\tau) = \varepsilon\cdot\sigma\cdot\sqrt{2b^2-1}\cdot\sin(nq + mp + \tau) - \varepsilon\cdot\sin(q + p + \tau). \tag{9.13}$$

We have taken $m \neq n$ to avoid that V depends only on a single variable $q + p$. The role of b (and σ) is to permit two different coupling constants. Then the control term is a sum of only five waves:

$$\begin{aligned} f(V) = \hat{\varepsilon}^2\cdot\sigma\cdot\sin((n-1)q + (m-1)p)\\ + \tilde{\varepsilon}^3\cdot[\sqrt{2b^2-1}\cdot(\sin(q+p+\tau) + \sin((2n-1)q + (2m-1)p + \tau))\\ - \sigma\cdot(\sin(nq + mp + \tau) + \sin((2-n)q + (2-m)p + \tau))]\end{aligned} \tag{9.14}$$

where

$$\hat{\varepsilon} := \frac{(2b^2-1)^{1/4}}{b}\cdot\left(\frac{1-\cos(\varepsilon\cdot b\cdot(m-n))}{|m-n|}\right)^{1/2} \approx |\varepsilon|\cdot|m-n|^{1/2}\cdot\left(\frac{2b^2-1}{4}\right)^{1/4} \tag{9.15}$$

$$\tilde{\varepsilon} := \frac{(4b^2-2)^{1/6}}{b}\cdot\left(\varepsilon\cdot b - \frac{\sin(\varepsilon\cdot b\cdot(m-n))}{m-n}\right)^{1/3} \approx \varepsilon\cdot|m-n|^{2/3}\cdot\left(\frac{2b^2-1}{18}\right)^{1/6}. \tag{9.16}$$

We have indicated the first order of the expansion of $\hat{\varepsilon}$ or $\tilde{\varepsilon}$ when ε is small, so that $f(V) = \mathcal{O}(V^2)$.

See also [4, 5] for some numerical experiments that prove the effectiveness of this method, when the coefficients $\varepsilon_{n,m,k}$ are taken to reflect some properties of a realistic field, in a Tokamak:

$$\varepsilon_{n,m,k} = \frac{\varepsilon}{(n^2+m^2)^{\frac{3}{2}}}\cdot\chi(1 \leqslant n^2 + m^2 \leqslant N^2)\cdot\chi(k=1) \tag{9.17}$$

for some constant ε proportional to the inverse of the (strong) magnetic field and for some 'cut-off' N. In that case the first term of the control, $\varepsilon^{(2)}_{n,m,k}$ is vanishing when $k \neq 0$. In [5], we also give some quantitative values of the parameters for the rigorous applicability of this control.

Let us summarize the method of control, in this case, where $\mathcal{R}V = 0$ and $\{H\} = \partial_\tau$:

$$\forall t \in \mathbb{R} \qquad \mathrm{e}^{t\{H+V+f(V)\}} = \mathrm{e}^{-\{\Gamma V\}} \cdot \mathrm{e}^{t\partial_\tau} \cdot \mathrm{e}^{\{\Gamma V\}} \tag{9.18}$$

so that the distance between the dynamical variable p (or q) at time t and at the initial time is

$$p_t - p_0 = (\mathrm{e}^{t\{H+V+f(V)\}} - 1)p_0. \tag{9.19}$$

Let us write p instead of p_0, and use the 'telescopic' formula:

$$\forall a, b, c: \qquad a \cdot b \cdot c - 1 = (a-1) + a \cdot (b-1) + a \cdot b \cdot (c-1). \tag{9.20}$$

And we replace the flow in (9.19) by its decomposition (9.18):

$$p_t - p = [(\mathrm{e}^{-\{\Gamma V\}} - 1) + \mathrm{e}^{-\{\Gamma V\}} \cdot (\mathrm{e}^{t\partial_\tau} - 1) + \mathrm{e}^{-\{\Gamma V\}} \cdot \mathrm{e}^{t\partial_\tau} \cdot (\mathrm{e}^{\{\Gamma V\}} - 1)]p. \tag{9.21}$$

But the middle term vanishes since $\partial_\tau p = 0$, so that

$$p_t - p = (\mathrm{e}^{-\{\Gamma V\}} - 1)p + \mathrm{e}^{-\{\Gamma V\}} \cdot \mathrm{e}^{t\partial_\tau} \cdot (\mathrm{e}^{\{\Gamma V\}} - 1)p. \tag{9.22}$$

And we can divide and multiply by $\{\Gamma V\}$, and use the antisymmetry (1.2):

$$p_t - p = \left(\frac{1 - \mathrm{e}^{-\{\Gamma V\}}}{\{\Gamma V\}}\right)\{p\}\Gamma V - \mathrm{e}^{-\{\Gamma V\}} \cdot \mathrm{e}^{t\partial_\tau} \cdot \left(\frac{\mathrm{e}^{\{\Gamma V\}} - 1}{\{\Gamma V\}}\right)\{p\}\Gamma V. \tag{9.23}$$

Finally, let us note that $\{p\} = -\partial_q$, i.e. formula (9.23) can be explicitly computed.

When we apply an approximate control term φ instead of the exact one $f(V)$, formula (9.18) becomes

$$\mathrm{e}^{t\{H+V+\varphi\}} = \mathrm{e}^{-\{\Gamma V\}} \cdot \mathrm{e}^{-\{\Gamma W\}} \cdot \mathrm{e}^{t\partial_\tau} \cdot \mathrm{e}^{t\{\mathcal{R}W\}} \cdot \mathrm{e}^{\{\Gamma W\}} \cdot \mathrm{e}^{\{\Gamma V\}} \tag{9.24}$$

where

$$W := F^{-1}(\mathrm{e}^{\{\Gamma V\}} \cdot (\varphi - f(V))). \tag{9.25}$$

When we take $\varphi = 0$, (9.25) is the beginning of the KAM recursive method. See ([3]) for some extensions and numerical tests of this theory for some dynamical systems. See also [7] for more details, in the case of quantum mechanics.

10. Example 2: quantum adiabatic transformation

An example of typical norm on $\mathcal{A}$ is given by an arbitrary 'weight' function $g(A, \Delta) > 0$. Let us define a norm on the Lie-algebra $\mathcal{A}$ by

$$\|V\| := \sup_A \sum_\Delta |V(A, \Delta)|/g(A, \Delta). \tag{10.1}$$

In the quantum mechanical case, we will take the usual L^2-operator norm on each 'block' $V(A, \Delta)$. This choice is irrelevant in the case where all the projectors P_A are finite-dimensional. We could have chosen $\sup_{A,\Delta}$ but this would have just been (approximately) a modification of the 'weight' function g.

Lemma 1. *We have*

$$\|\mathcal{R}V\| \leqslant \|V\| \qquad \text{and} \qquad \|\mathcal{N}V\| \leqslant \|V\|. \tag{10.2}$$

Proof. The operators $\mathcal{R}$ and $\mathcal{N}$ are implemented by some characteristic functions, cf (1.14), (1.15), (1.26) and (1.27). So the sup and sum in (10.1) are restricted by some conditions, therefore the norm decreases. □

Hence hypothesis 3 is fulfilled. Let us consider as before a Hamiltonian H but where now the perturbation depends on time. We need to extend the algebra and the full Hamiltonian is now $D+H$ where D is the 'derivative with respect to time'. We can apply proposition 1 (5.1) but we still use the same operator Γ i.e. the pseudo-inverse of the bracket with H, and also the same function F defined in (2.1). Let us rewrite (5.1):

$$\forall\, W \in \mathcal{A} \qquad H + F(W) = \mathrm{e}^{-\{\Gamma W\}}(H + \mathcal{R}W). \tag{10.3}$$

Now we want some information on $D + H + V$:

$$D + H + F(W) = D + \mathrm{e}^{-\{\Gamma W\}}(H + \mathcal{R}W) = \mathrm{e}^{-\{\Gamma W\}}(\mathrm{e}^{\{\Gamma W\}} D + H + \mathcal{R}W) \tag{10.4}$$

where

$$W := F^{-1}(V). \tag{10.5}$$

Hence

$$D + H + F(W) = \mathrm{e}^{-\{\Gamma W\}}\left(D + \frac{\mathrm{e}^{\{\Gamma W\}} - 1}{\{\Gamma W\}} \cdot \{\Gamma W\} D + H + \mathcal{R}W\right) \tag{10.6}$$

so, using (1.2), and the notation $\dot{W} := \{D\}W$

$$D + H + V = \mathrm{e}^{-\{\Gamma W\}}(D + H_1 + V_1) \tag{10.7}$$

with

$$H_1 := H + \mathcal{R}W \qquad V_1 := -\frac{\mathrm{e}^{\{\Gamma W\}} - 1}{\{\Gamma W\}} \cdot \Gamma \dot{W}. \tag{10.8}$$

Indeed $\{D\}\Gamma W = \Gamma\{D\}W = \Gamma\dot{W}$. Formula (10.7) is useful if the derivative $\dot{V}$ of the perturbation with respect to time is 'smaller' than the perturbation V itself. This is the 'adiabatic hypothesis'. In that case we can also show that $\dot{W}$ is small, since $W \approx V$. So the new perturbation is 'smaller' than the original perturbation, or more qualitatively it is approximately $\Gamma\dot{W}$ i.e. $\{H\}^{-1}\{D\}V$. This means that it has lost some regularity in time, but gained some regularity in its 'spatial' dependence, since generally $\{H\}^{-1}$ is a regularizing operator. The difference between H and the new H_1 is of the same size as the original perturbation, but 'diagonalized' with respect to H since it commutes with it, cf also [6] for an iterative proof of this result.

Let us finally note that it is possible to iterate this procedure:

$$D + H + V = \mathrm{e}^{-\{\Gamma W\}}\,\mathrm{e}^{-\{\Gamma_1 W_1\}}(D + H_2 + V_2) \qquad \text{where} \quad W_1 := F_1^{-1}(V_1) \tag{10.9}$$

and

$$H_2 := H_1 + \mathcal{R}_1 W_1 \qquad V_2 := -\frac{\mathrm{e}^{\{\Gamma_1 W_1\}} - 1}{\{\Gamma_1 W_1\}} \cdot \Gamma_1 \dot{W}_1 \tag{10.10}$$

and with $\mathcal{R}_1$, $\mathcal{N}_1$, Γ_1, F_1 defined as before, but with H replaced by H_1, and so on. Generally this iteration does not converge: so we must stop it at an optimized order. This has been done in [10], with a different method and framework.

Finally let us explicit the norm (7.1) of the operator Γ for the norm chosen in (10.1):

$$|||\Gamma||| \leqslant \sup_{A, A' \text{s.t.} A \neq A'} \frac{\psi(A, A')}{|h(A) - h(A')|} \tag{10.11}$$

where

$$\psi(A, A') := g_{A,A'} \cdot \sup_{A''} \max\left(\frac{g_{A,A''}}{g_{A',A''}}, \frac{g_{A',A''}}{g_{A,A''}}\right) \tag{10.12}$$

and $g_{A,A'} = g(A, A'-A)$. The proof is an easy estimation of the norm (7.1). Hence hypothesis 2 (7.2) requires that

$$\forall\, A, A' \quad \text{s.t. } A \neq A' \qquad |h(A) - h(A')| \geqslant \gamma \cdot \psi(A, A') \tag{10.13}$$

for some constant $\gamma > 0$. This is a 'diophantine' condition. It depends on the choice of the weight function g, i.e. on the regularity of the perturbations V we want to consider. For instance, when $\mathcal{H} = \mathbb{N}^*$

$$\text{if} \quad g_{A,A'} = \varphi_A \cdot \varphi_{A'} \qquad \text{then} \quad \psi(A, A') = \max\left(\varphi_A^2, \varphi_{A'}^2\right). \tag{10.14}$$

As a particular case we can take, for some $\alpha \geqslant 1$:

$$h(A) = A^\alpha \qquad \text{and} \qquad \varphi_A = A^{\frac{\alpha-1}{2}} \tag{10.15}$$

for which condition (10.13) is satisfied.

See [7] for more details.

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