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A non self-adjoint model on a two dimensional noncommutative space with unbound metric

Fabio Bagarello[◦] and Andreas Fring[•]

[◦] *Dipartimento di Energia, Ingegneria dell'Informazione e Modelli Matematici, Facoltà di Ingegneria, Università di Palermo, I-90128 Palermo, Italy and INFN, Sezione di Torino, Italy*

[•] *Department of Mathematical Science, City University London, Northampton Square, London EC1V 0HB, UK*

E-mail: fabio.bagarello@unipa.it, a.fring@city.ac.uk

ABSTRACT: We demonstrate that a non self-adjoint Hamiltonian of harmonic oscillator type defined on a two-dimensional noncommutative space can be diagonalized exactly by making use of pseudo-bosonic operators. The model admits an antilinear symmetry and is of the type studied in the context of \mathcal{PT} -symmetric quantum mechanics. Its eigenvalues are computed to be real for the entire range of the coupling constants and the biorthogonal sets of eigenstates for the Hamiltonian and its adjoint are explicitly constructed. We show that despite the fact that these sets are complete and biorthogonal, they involve an unbounded metric operator and therefore do not constitute (Riesz) bases for the Hilbert space $\mathcal{L}^2(\mathbb{R}^2)$, but instead only D-quasi bases. As recently proved by one of us (FB), this is sufficient to deduce several interesting consequences.

1. Introduction

In the last 15 years more and more of physicists and mathematicians have developed an interest in non-Hermitian and non self-adjoint operators possessing real eigenvalues. Such type of models have been investigated before, but the more recent interest has been initiated by the seminal paper [1] in which the complex cubic potential and its close relatives have been studied. The original considerations, focussing mainly on the aspect of the possibility to formulate consistent quantum mechanical systems, have broadened quickly and are partly replaced by a more general analysis of related aspects. Many experiments [2, 3, 4, 5] have now been carried out, mainly for optical analogues to the quantum mechanical systems, exploiting \mathcal{PT} -symmetric phase transitions where real eigenvalues merge into two complex conjugate pairs, to obtain gain and loss structures. We refer the reader to [6, 7, 8] for some reviews on what is commonly named *quasi-Hermitian* [9, 10], *pseudo-Hermitian* [11, 12] or \mathcal{PT} -symmetric [13, 1] quantum mechanics. However, it was recently pointed

out by Krejcirik and Siegl [14] that more mathematically oriented treatments of these type of Hamiltonians are required, as for instance the complex cubic potential lacks to possess a Riesz basis of eigenstates. Therefore we can still not associate a standard quantum mechanical interpretation to this model. The purpose of this paper is to shed more light on these issues.

Modifying recent ideas [15], one of us has newly introduced the notion of \mathcal{D} -pseudo bosons (\mathcal{D} -pbs), [16], and used them in connection with several physical systems, whose Hamiltonians are non self-adjoint operators, [17]. Among other aspects, it was shown that \mathcal{D} -pbs could be useful in the analysis of a two-dimensional harmonic oscillator described by the Hamiltonian

$$\hat{H} = \frac{1}{2}(\hat{p}_1^2 + \hat{x}_1^2) + \frac{1}{2}(\hat{p}_2^2 + \hat{x}_2^2) + i[A(\hat{x}_1 + \hat{x}_2) + B(\hat{p}_1 + \hat{p}_2)], \quad (1.1)$$

where (\hat{x}_j, \hat{p}_j) are noncommutative operators satisfying $[\hat{x}_j, \hat{p}_k] = i\delta_{j,k}\mathbb{1}$, $[\hat{x}_j, \hat{x}_k] = i\theta\epsilon_{j,k}\mathbb{1}$, $[\hat{p}_j, \hat{p}_k] = i\tilde{\theta}\epsilon_{j,k}\mathbb{1}$, where θ and $\tilde{\theta}$ are two real small parameters, measuring the non commutativity of the system. In [17] a perturbative expansion in θ and $\tilde{\theta}$ was set up and it was shown, in particular, that if one neglects all the terms which are at least quadratic in θ and $\tilde{\theta}$ we can construct explicitly the eigenvectors of (the approximated version of) \hat{H} and deduce the related eigenvalues.

In this paper we show that, if the non commutativity is restricted to the spatial variables only, i.e. if $\tilde{\theta} = 0$, then \hat{H} , and a slightly generalized version of it, can be exactly diagonalized in terms of \mathcal{D} -pbs. The corresponding eigenbases are biorthonormal, but involve a metric operator that is unbounded, together with its inverse. Thus we will draw a similar conclusion as reached in [14] and, more recently, in [18].

It may be worth to underline that these results, all together, suggest that several common believes usually taken for granted in the physical literature on these topics require some more care than usually adopted. For instance, in [19] (as well as in many other papers, [20]), the biorthogonal sets of eigenstates of a rather general H , with $H^\dagger \neq H$, are used to produce a resolution of the identity. In other words, they are used as bases in the Hilbert space. However, the results in [14, 18], and those given in this paper, show that this is not always possible, even for extremely simple models. This, we believe, helps clarifying the situation, showing that many claims need to be analyzed in more details.

This article is organized as follows: in the next section we review the definition and a few central results on \mathcal{D} -pbs. In section 3 we introduce the 2d-harmonic oscillator with linear term in the momenta and position on a noncommutative flat space and we analyze it in terms of \mathcal{D} -pbs. We provide the computation of how it may be written in terms of \mathcal{D} -pb number operators and subsequently we verify the underlying assumptions, needed to have something more than just a formal theory. This will allow for the construction of biorthonormal sets, which are, however, shown not to be Riesz bases and not even bases, but just \mathcal{D} -quasi bases. Our conclusions are stated in section 4.

2. Pseudo-bosons, generalities

We briefly review here few definitions and central properties of \mathcal{D} -pbs. More details can

be found in [16].

Let \mathcal{H} be a given Hilbert space with scalar product $\langle \cdot, \cdot \rangle$ and related norm $\|\cdot\|$. Furthermore, let a and b be two operators acting on \mathcal{H} , with domains $D(a)$ and $D(b)$ respectively, a^\dagger and b^\dagger their respective adjoints, and let \mathcal{D} be a dense subspace of \mathcal{H} such that $a^\sharp \mathcal{D} \subseteq \mathcal{D}$ and $b^\sharp \mathcal{D} \subseteq \mathcal{D}$, where x^\sharp is x or x^\dagger . It is worth noticing that we are not requiring here that \mathcal{D} coincides with either $D(a)$ or $D(b)$. Nevertheless, for obvious reasons, $\mathcal{D} \subseteq D(a^\sharp)$ and $\mathcal{D} \subseteq D(b^\sharp)$.

Definition: The operators (a, b) are \mathcal{D} -pseudo-bosonic if, for all $f \in \mathcal{D}$, we have

$$abf - baf = f. \quad (2.1)$$

Sometimes, to simplify the notation, instead of (2.1) we will simply write $[a, b] = \mathbb{1}$, having in mind that both sides of this equation have to act on $f \in \mathcal{D}$.

Our working assumptions are the following:

Assumption \mathcal{D} -pb 1: There exists a non-zero $\varphi_0 \in \mathcal{D}$ such that $a\varphi_0 = 0$.

Assumption \mathcal{D} -pb 2: There exists a non-zero $\Psi_0 \in \mathcal{D}$ such that $b^\dagger\Psi_0 = 0$.

Then, if (a, b) satisfy the above definition, it is obvious that $\varphi_0 \in D^\infty(b)$ and that $\Psi_0 \in D^\infty(a^\dagger)$, with $D^\infty(x)$ denoting the common domain of all powers of x . Thus we can define the following vectors, all belonging to \mathcal{D} :

$$\varphi_n := \frac{1}{\sqrt{n!}} b^n \varphi_0, \quad \Psi_n := \frac{1}{\sqrt{n!}} a^{\dagger n} \Psi_0, \quad (2.2)$$

for $n \geq 0$. As in [16] we introduce the sets $\mathcal{F}_\Psi = \{\Psi_n, n \geq 0\}$ and $\mathcal{F}_\varphi = \{\varphi_n, n \geq 0\}$. Once again, since \mathcal{D} is stable under the action of a^\sharp and b^\sharp , we deduce that each φ_n and each Ψ_n belongs to the domains of a^\sharp , b^\sharp and N^\sharp , where $N := ba$.

It is now straightforward to deduce the following lowering and raising relations:

$$\begin{aligned} a\varphi_n &= \sqrt{n} \varphi_{n-1}, \quad a\varphi_0 = 0, & b^\dagger\Psi_n &= \sqrt{n} \Psi_{n-1}, \quad b^\dagger\Psi_0 = 0, & \text{for } n \geq 1, \\ a^\dagger\Psi_n &= \sqrt{n+1} \Psi_{n+1}, & b\varphi_n &= \sqrt{n+1} \varphi_{n+1}, & \text{for } n \geq 0, \end{aligned} \quad (2.3)$$

as well as the following eigenvalue equations: $N\varphi_n = n\varphi_n$ and $N^\dagger\Psi_n = n\Psi_n$ for $n \geq 0$. As a consequence of these equations, choosing the normalization of φ_0 and Ψ_0 in such a way $\langle \varphi_0, \Psi_0 \rangle = 1$, we deduce that

$$\langle \varphi_n, \Psi_m \rangle = \delta_{n,m}, \quad (2.4)$$

for all $n, m \geq 0$. The third assumption originally introduced in [16] is the following:

Assumption \mathcal{D} -pb 3: \mathcal{F}_φ is a basis for \mathcal{H} .

This is equivalent to the request that \mathcal{F}_Ψ is a basis for \mathcal{H} as well, [16]. In particular, if \mathcal{F}_φ and \mathcal{F}_Ψ are Riesz bases for \mathcal{H} , the \mathcal{D} -pbs were called *regular*.

In [16] also a weaker version of Assumption \mathcal{D} -pb 3 has been introduced, useful for concrete physical applications: for that, let \mathcal{G} be a suitable dense subspace of \mathcal{H} . Two

biorthogonal sets $\mathcal{F}_\eta = \{\eta_n \in \mathcal{G}, g \geq 0\}$ and $\mathcal{F}_\Phi = \{\Phi_n \in \mathcal{G}, g \geq 0\}$ were called \mathcal{G} -quasi bases if, for all $f, g \in \mathcal{G}$, the following holds:

$$\langle f, g \rangle = \sum_{n \geq 0} \langle f, \eta_n \rangle \langle \Phi_n, g \rangle = \sum_{n \geq 0} \langle f, \Phi_n \rangle \langle \eta_n, g \rangle. \quad (2.5)$$

It is clear that, while Assumption \mathcal{D} -pb 3 implies (2.5), the reverse is false. However, if \mathcal{F}_η and \mathcal{F}_Φ satisfy (2.5), we still have some (weak) form of resolution of the identity. Now Assumption \mathcal{D} -pb 3 is replaced by the following:

Assumption \mathcal{D} -pbw 3: \mathcal{F}_φ and \mathcal{F}_Ψ are \mathcal{G} -quasi bases.

Let now assume that Assumption \mathcal{D} -pb 1, \mathcal{D} -pb 2, and \mathcal{D} -pbw 3 are satisfied, with $\mathcal{G} = \mathcal{D}$, and let us consider a self-adjoint, invertible, operator Θ , which leaves, together with Θ^{-1} , \mathcal{D} invariant: $\Theta\mathcal{D} \subseteq \mathcal{D}$, $\Theta^{-1}\mathcal{D} \subseteq \mathcal{D}$. Then, as in [16], we say that (a, b^\dagger) are Θ -conjugate if $af = \Theta^{-1}b^\dagger\Theta f$, for all $f \in \mathcal{D}$. Moreover, we can check that, for instance, (a, b^\dagger) are Θ -conjugate if and only if (b, a^\dagger) are Θ -conjugate and that, assuming that $\langle \varphi_0, \Theta\varphi_0 \rangle = 1$, (a, b^\dagger) are Θ -conjugate if and only if $\Psi_n = \Theta\varphi_n$, for all $n \geq 0$. Finally, if (a, b^\dagger) are Θ -conjugate, then $\langle f, \Theta f \rangle > 0$ for all non zero $f \in \mathcal{D}$. The details of these proofs can be found in [18]. Notice also that, not surprisingly, we also deduce that $Nf = \Theta^{-1}N^\dagger\Theta f$, for all $f \in \mathcal{D}$.

3. Noncommutative two dimensional harmonic oscillator with linear terms

Let us now consider the non self-adjoint two dimensional harmonic oscillator with linear terms in the momenta and positions

$$\tilde{H} = \frac{1}{2m}(\tilde{p}_1^2 + \tilde{p}_2^2) + \frac{m\omega^2}{2}(\tilde{x}_1^2 + \tilde{x}_2^2) + i\alpha_1\tilde{x}_1 + \alpha_2\tilde{x}_2 + \alpha_3\tilde{p}_1 + i\alpha_4\tilde{p}_2, \quad (3.1)$$

on the noncommutative flat space with the nonvanishing commutators $[\tilde{x}_1, \tilde{x}_2] = i\theta$, $[\tilde{x}_j, \tilde{p}_j] = i\hbar$ for $j = 1, 2$. Here θ and α_i for $i = 1, 2, 3, 4$ are real dimensionful parameters. Note that this Hamiltonian is non self-adjoint even when viewed on a standard space. However, \tilde{H} is constructed in such a way that it is left invariant with respect to the antilinear symmetry \mathcal{PT}_- : $\tilde{x}_1 \rightarrow -\tilde{x}_1$, $\tilde{x}_2 \rightarrow \tilde{x}_2$, $\tilde{p}_1 \rightarrow \tilde{p}_1$, $\tilde{p}_2 \rightarrow -\tilde{p}_2$ and $i \rightarrow -i$ [21]. Thus in the general spirit of \mathcal{PT} -symmetric quantum mechanics [13, 1] the Hamiltonian is guaranteed to have real eigenvalues provided that its eigenfunctions are eigenstates of \mathcal{PT}_- . Evidently in atomic units, $m = \omega = \hbar = 1$, \tilde{H} reduces to \hat{H} for $\alpha_1 \rightarrow A$, $\alpha_2 \rightarrow -iA$, $\alpha_3 \rightarrow iB$ and $\alpha_4 \rightarrow B$. We also notice that \mathcal{PT}_- is no longer a symmetry of \tilde{H} , i.e. $[\mathcal{PT}_-, \tilde{H}] \neq 0$.

Our aim here is to employ \mathcal{D} -pbs to diagonalize \tilde{H} exactly, instead of using a perturbative approach as in [17, 22] and to determine its spectrum. For this purpose we convert the Hamiltonian first from a flat noncommutative space to one in terms of standard canonical variables x_i and p_i for $i = 1, 2$ satisfying the canonical commutation relations $[x_j, p_j] = i\hbar$ and $[x_i, x_j] = [p_i, p_j] = 0$. This is achieved by a standard Bopp shift $\tilde{x}_1 \rightarrow x_1 - \frac{\theta}{2\hbar}p_2$,

$\tilde{x}_2 \rightarrow x_2 + \frac{\theta}{2\hbar}p_1$, $p_1 \rightarrow p_1$ and $p_2 \rightarrow p_2$. The Hamiltonian in (3.1) then acquires the form

$$\begin{aligned} \tilde{H} = & \left(\frac{1}{2m} + \frac{m\omega^2\theta^2}{8\hbar^2} \right) (p_1^2 + p_2^2) + \frac{m\omega^2}{2}(x_1^2 + x_2^2) + \frac{m\omega^2\theta}{2\hbar}(x_2p_1 - x_1p_2) \\ & + i\alpha_1x_1 + \alpha_2x_2 + \left(\alpha_3 + \frac{\alpha_2\theta}{2\hbar} \right) p_1 + i \left(\alpha_4 - \frac{\alpha_1\theta}{2\hbar} \right) p_2. \end{aligned} \quad (3.2)$$

We now attempt to re-express this Hamiltonian in terms of pseudo-bosonic number operators $N_i = b_i a_i$ as

$$\tilde{H} = \gamma_1 N_1 + \gamma_2 N_2 + \gamma_0 \quad \text{for } \gamma_0, \gamma_1, \gamma_2 \in \mathbb{R}, \quad (3.3)$$

where the operators a_i and b_i obey the two dimensional pseudo-bosonic commutation relations

$$[a_j, b_k] = i\delta_{jk}, \quad [a_j, a_k] = [b_j, b_k] = 0, \quad \text{for } j, k = 1, 2. \quad (3.4)$$

For this purpose we represent the pseudo-bosonic operators a_i and b_i in terms of standard bosonic creation and annihilation operators A_i^\dagger and A_i , respectively,

$$a_1 = \frac{1}{\sqrt{2}}(A_1 + iA_2) + i\beta_1, \quad b_1 = \frac{1}{\sqrt{2}}(A_1^\dagger - iA_2^\dagger) + i\beta_3, \quad (3.5)$$

$$a_2 = -\frac{1}{\sqrt{2}}(iA_1 + A_2) + \beta_2, \quad b_2 = \frac{1}{\sqrt{2}}(iA_1^\dagger - A_2^\dagger) + \beta_4, \quad (3.6)$$

with $[A_j, A_k^\dagger] = i\delta_{jk}$, $[A_j, A_k] = [A_j^\dagger, A_k^\dagger] = 0$ for $j, k = 1, 2$ and $\beta_i \in \mathbb{C}$ for $i = 1, 2, 3, 4$. Furthermore we represent the A_i^\dagger and A_i in terms of the standard canonical variables

$$A_1 = \sqrt{\frac{M\omega}{2\hbar}}x_1 + i\sqrt{\frac{1}{2\hbar M\omega}}p_1, \quad A_2 = \sqrt{\frac{M\omega}{2\hbar}}x_2 + i\sqrt{\frac{1}{2\hbar M\omega}}p_2, \quad (3.7)$$

$$A_1^\dagger = \sqrt{\frac{M\omega}{2\hbar}}x_1 - i\sqrt{\frac{1}{2\hbar M\omega}}p_1, \quad A_2^\dagger = \sqrt{\frac{M\omega}{2\hbar}}x_2 - i\sqrt{\frac{1}{2\hbar M\omega}}p_2. \quad (3.8)$$

We note that the pseudo-bosonic operators reduce to standard boson operators with $b_i = a_i^\dagger$ if and only if for $\beta_1 = -\bar{\beta}_3$ and $\beta_2 = \bar{\beta}_4$. Upon substitution we compare now (3.3) and (3.2), which become identical subject to the constraints

$$\beta_1 = \frac{\Omega(\alpha_1 + \alpha_2) + 2\hbar m\omega(\alpha_3 - \alpha_4)}{(\Omega + \theta m\omega)\sqrt{2m\Omega\omega^3}}, \quad \beta_2 = \frac{\Omega(\alpha_1 - \alpha_2) + 2\hbar m\omega(\alpha_3 + \alpha_4)}{(\Omega - \theta m\omega)\sqrt{2m\Omega\omega^3}}, \quad (3.9)$$

$$\beta_3 = \frac{\Omega(\alpha_1 - \alpha_2) - 2\hbar m\omega(\alpha_3 + \alpha_4)}{(\Omega + \theta m\omega)\sqrt{2m\Omega\omega^3}}, \quad \beta_4 = \frac{-\Omega(\alpha_1 + \alpha_2) + 2\hbar m\omega(\alpha_3 - \alpha_4)}{(\Omega - \theta m\omega)\sqrt{2m\Omega\omega^3}}, \quad (3.10)$$

$$\gamma_0 = \frac{1}{2}\omega [\Omega (1 + \beta_1\beta_3 - \beta_2\beta_4) + \theta m\omega (\beta_1\beta_3 + \beta_2\beta_4)], \quad (3.11)$$

$$\gamma_1 = \frac{1}{2}\omega (\Omega + \theta m\omega), \quad \gamma_2 = \frac{1}{2}\omega (\Omega - \theta m\omega), \quad M = \frac{2m\hbar}{\Omega}, \quad (3.12)$$

where $\Omega := \sqrt{4\hbar^2 + \theta^2 m^2 \omega^2}$. If we are now able to construct eigenstates $\Psi_{\underline{n}}$ for the pseudo-bosonic number operators such that $N_i \varphi_{\underline{n}} = \hbar \omega n_i \varphi_{\underline{n}}$, the eigenvalues for \tilde{H} are immediately computed from (3.3) to

$$E_{n_1, n_2} = \gamma_1 \hbar \omega n_1 + \gamma_2 \hbar \omega n_2 + \gamma_0. \quad (3.13)$$

We observe from (3.9) to (3.12) that the constants $\gamma_i \in \mathbb{R}$ for $i = 0, 1, 2$ are real and consequently the energy E_{n_1, n_2} is also real. Furthermore, we observe that the presence of the linear terms in (3.1), that is $\alpha_i \neq 0$ for $i = 1, 2, 3, 4$, prevents us from using a standard bosonic oscillator algebra and we are forced to employ pseudo-bosons. This is seen from the fact that the pseudo-bosonic operator reduce to standard boson operators if and only if for $\beta_1 = -\bar{\beta}_3$ and $\beta_2 = \bar{\beta}_4$. However, our constraints (3.9) and (3.10) imply that in this boson case some linear terms in our Hamiltonian have to vanish, that is $\alpha_1 = \alpha_4 = 0$.

Furthermore we notice that for the reduction of \tilde{H} to \hat{H} for $\alpha_1 \rightarrow A$, $\alpha_2 \rightarrow iA$, $\alpha_3 \rightarrow iB$, $\alpha_4 \rightarrow B$ we obtain $\beta_1 = \bar{\beta}_3$ and $\beta_2 = -\bar{\beta}_4$, such that γ_0 and therefore E_{n_1, n_2} remain real. In this case the \mathcal{PT} -symmetry is broken and it remains unclear which antilinear symmetry, if any, is responsible for keeping the spectrum real.

Let us now verify that eigenstates $\varphi_{\underline{n}}$ and those of the adjoint of the Hamiltonian, $\Psi_{\underline{n}}$, are well defined, really exist and most crucially whether they constitute a Riesz basis, or even a basis.

3.1 Verification of the pseudo-bosonic assumptions

For simplicity let us now adopt atomic units. We commence by introducing the operators

$$\hat{a}_i := \lim_{\beta_i \rightarrow 0} a_i, \quad \hat{a}_i^\dagger := \lim_{\beta_i \rightarrow 0} b_i, \quad (3.14)$$

which, from (3.5)-(3.6), satisfy the standard bosonic canonical commutation relations, $[\hat{a}_i, \hat{a}_j^\dagger] = \delta_{i,j} \mathbb{1}$, $[\hat{a}_i, \hat{a}_j] = 0$, for $i, j = 1, 2$. Then, introducing the unitary operators

$$D_i(z) := \exp \left\{ \bar{z} \hat{a}_i - z \hat{a}_i^\dagger \right\}, \quad D(\underline{z}) := D_1(z_1) D_2(z_2), \quad (3.15)$$

we compute

$$a_i = \hat{a}_i + \nu_i = D(\underline{\nu}) \hat{a}_i D^{-1}(\underline{\nu}), \quad b_i = \hat{a}_i^\dagger + \mu_i = D(\underline{\mu}) \hat{a}_i^\dagger D^{-1}(\underline{\mu}), \quad (3.16)$$

for $i = 1, 2$ with $\underline{\nu} := \{i\beta_1, \beta_2\}$, $\underline{\mu} := \{-i\bar{\beta}_3, \bar{\beta}_4\}$. An orthonormal basis for $\mathcal{H} = \mathcal{L}^2(\mathbb{R}^2)$ is then constructed easily: Let $e_{0,0} = e_{\underline{0}}$ be the vacuum of \hat{a}_1 and \hat{a}_2 , that is $\hat{a}_i e_{\underline{0}} = 0$ for $i = 1, 2$. Then as common for the purely bosonic case, we introduce

$$e_{n_1, n_2} = e_{\underline{n}} := \frac{1}{\sqrt{n_1! n_2!}} (\hat{a}_1^\dagger)^{n_1} (\hat{a}_2^\dagger)^{n_2} e_{\underline{0}}, \quad (3.17)$$

and the related orthonormal basis $\mathcal{F}_e = \{e_{\underline{n}}, n_1, n_2 \geq 0\}$. Of course for the bosonic number operator $\hat{n}_i := \hat{a}_i^\dagger \hat{a}_i$ we have $\hat{n}_i e_{\underline{n}} = n_i e_{\underline{n}}$.

In order to verify the assumptions of section 2, we first seek to construct $\varphi_{\underline{0}}$, i.e. the vacuum of a_i satisfying $a_1 \varphi_{\underline{0}} = a_2 \varphi_{\underline{0}} = 0$. Evidently this holds if, and only if, $\hat{a}_i (D^{-1}(\underline{\nu}) \varphi_{\underline{0}}) = 0$ for $i = 1, 2$. This implies that $\varphi_{\underline{0}} = D(\underline{\nu}) e_{\underline{0}}$, up to a normalization which will be fixed below. Notice that, due to fact that $D(\underline{\nu})$ is unitary, and therefore everywhere defined, $\varphi_{\underline{0}}$ is well defined.

Similarly we derive $\Psi_{\underline{0}}$, the vacuum for b_i^\dagger . We require $b_1^\dagger \Psi_{\underline{0}} = b_2^\dagger \Psi_{\underline{0}} = 0$ which can be rewritten as $\hat{a}_i (D^{-1}(\underline{\mu}) \Psi_{\underline{0}}) = 0$ for $i = 1, 2$. These equations are solved by $\Psi_{\underline{0}} = N_\Psi D(\underline{\mu}) e_{\underline{0}}$,

which, due to the unitarity of $D(\underline{\mu})$ is again well defined. Here N_Ψ is a normalization needed to ensure the normalization $\langle \varphi_{\underline{0}}, \Psi_{\underline{0}} \rangle = 1$. It is computed to

$$N_\Psi^2 = \frac{\langle \varphi_{\underline{0}}, \varphi_{\underline{0}} \rangle}{\langle \Psi_{\underline{0}}, \Psi_{\underline{0}} \rangle} = \exp \left[|\beta_1|^2 + |\beta_2|^2 - |\beta_3|^2 - |\beta_4|^2 - 2 \operatorname{Re}(\beta_1 \beta_2) - 2 \operatorname{Re}(\beta_3 \beta_4) \right]. \quad (3.18)$$

Evidently for $\beta_2 = -\beta_3$ and $\beta_1 = \beta_4$ this reduces to the standard bosonic normalization, as is expected.

Remark: These results could have also been found quite easily by solving the equations directly in the coordinate representation. For instance, $a_1 \varphi_{\underline{0}} = a_2 \varphi_{\underline{0}} = 0$ are equivalent to the differential equations

$$(x_1 + \partial_{x_1} + ix_2 + i \partial_{x_2} + 2i\beta_1) \varphi_{\underline{0}}(x_1, x_2) = (-ix_1 - i \partial_{x_1} - x_2 - \partial_{x_2} + 2\beta_2) \varphi_{\underline{0}}(x_1, x_2) = 0, \quad (3.19)$$

solved by $\varphi_{\underline{0}}(x_1, x_2) \propto e^{-\frac{1}{2}(x_1^2+x_2^2)-i(\beta_1+\beta_2)x_1-(\beta_1-\beta_2)x_2}$. Similarly we find $\Psi_{\underline{0}}(x_1, x_2) \propto e^{-\frac{1}{2}(x_1^2+x_2^2)+i(\beta_3-\beta_4)x_1+(\beta_3+\beta_4)x_2}$. We see that both of these functions belong, for instance, to the set $\mathcal{S}(\mathbb{R}^2)$ of C^∞ -functions which, together with their derivatives, decrease faster to zero than any inverse power of x_1 and x_2 . However, this property might not be enough for our purposes, since as we have outlined in section 2, we need to identify a set \mathcal{D} , dense in \mathcal{H} , which not only contains $\varphi_{\underline{0}}$ and $\Psi_{\underline{0}}$, but which is in addition also stable under the action of a_j^\sharp , b_j^\sharp , and other relevant operators. It is convenient to introduce, therefore, the following set:

$$\mathcal{D} = \left\{ f(x_1, x_2) \in \mathcal{S}(\mathbb{R}^2), \text{ such that } e^{k_1 x_1 + k_2 x_2} f(x_1, x_2) \in \mathcal{S}(\mathbb{R}^2), \forall k_1, k_2 \in \mathbb{C} \right\}. \quad (3.20)$$

\mathcal{D} is dense in \mathcal{H} , since it contains the set $D(\mathbb{R}^2)$ of the C^∞ -functions with compact support.

Following section 2, we are now interested in deducing the properties of the vectors $\varphi_{\underline{n}} = \frac{1}{\sqrt{n_1! n_2!}} b_1^{n_1} b_2^{n_2} \varphi_{\underline{0}}$ and $\Psi_{\underline{n}} = \frac{1}{\sqrt{n_1! n_2!}} (a_1^\dagger)^{n_1} (a_2^\dagger)^{n_2} \Psi_{\underline{0}}$. We notice that both $\varphi_{\underline{n}}$ and $\Psi_{\underline{n}}$ necessarily belong to \mathcal{D} for all \underline{n} , because of the stability of \mathcal{D} under the action of b_i and a_i^\dagger , and the previously established fact that $\varphi_{\underline{0}}, \Psi_{\underline{0}} \in \mathcal{D}$. The formulae (3.16) state how the pseudo-bosonic operators (a_i, b_i) are related to the bosonic operators $(\hat{a}_i, \hat{a}_i^\dagger)$ by means of the in general two different unitary operators $D(\underline{\nu})$ and $D(\underline{\mu})$.

A single operator could be used if we introduce the operators

$$V_i(z, w) := \exp \left\{ \bar{w} \hat{a}_i - z \hat{a}_i^\dagger \right\}, \quad V(\underline{\nu}, \underline{\mu}) := V_1(\nu_1, \mu_1) V_2(\nu_2, \mu_2). \quad (3.21)$$

Now we compute

$$a_i = V(\underline{\nu}, \underline{\mu}) \hat{a}_i V^{-1}(\underline{\nu}, \underline{\mu}), \quad b_i = V(\underline{\nu}, \underline{\mu}) \hat{a}_i^\dagger V^{-1}(\underline{\nu}, \underline{\mu}), \quad (3.22)$$

which, in contrast to (3.16), only involve a single, albeit in general unbounded, operator to relate the (a_i, b_i) to the $(\hat{a}_i, \hat{a}_i^\dagger)$. We also check directly

$$a_i^\dagger = V(\underline{\mu}, \underline{\nu}) \hat{a}_i V^{-1}(\underline{\mu}, \underline{\nu}), \quad b_i^\dagger = V(\underline{\mu}, \underline{\nu}) \hat{a}_i V^{-1}(\underline{\mu}, \underline{\nu}). \quad (3.23)$$

A immediate consequence of these formulae are the following relations between the various number operators: $\hat{n}_i = V^{-1}(\underline{\nu}, \underline{\mu})N_iV(\underline{\nu}, \underline{\mu}) = V^{-1}(\underline{\mu}, \underline{\nu})N_i^\dagger V(\underline{\mu}, \underline{\nu})$, which in turns implies that

$$N_i = T(\underline{\nu}, \underline{\mu})N_i^\dagger T^{-1}(\underline{\nu}, \underline{\mu}), \quad (3.24)$$

where $T(\underline{\nu}, \underline{\mu}) := V(\underline{\nu}, \underline{\mu})V^{-1}(\underline{\mu}, \underline{\nu})$. Needless to say, all these equalities and definitions are well defined on \mathcal{D} , but not on the whole \mathcal{H}^1 . Incidentally we observe that $T(\underline{\nu}, \underline{\nu}) = \mathbb{1}$. This is in agreement with the fact that, when $\underline{\mu} = \underline{\nu}$, the operator $V(\underline{\nu}, \underline{\mu})$ is bounded with bounded inverse, see below.

By a similar reasoning as above applied for the construction of the vacuum state we now deduce that

$$\varphi_{\underline{n}} = V(\underline{\nu}, \underline{\mu})e_{\underline{n}}, \quad \Psi_{\underline{n}} = N_{\Psi}V(\underline{\mu}, \underline{\nu})e_{\underline{n}}. \quad (3.25)$$

In analogy with [18], we see that, while $V(\underline{\nu}, \underline{\nu}) = D(\underline{\nu})$ is a unitary operator and as a consequence bounded, the operator $V(\underline{\nu}, \underline{\mu})$, as well as its inverse, is unbounded for $\underline{\nu} \neq \underline{\mu}$. The crucial conclusion from this is that the two sets $\mathcal{F}_{\varphi} = \{\varphi_{\underline{n}}\}$ and $\mathcal{F}_{\Psi} = \{\Psi_{\underline{n}}\}$ cannot be Riesz bases. In fact, they are both related to the orthonormal basis \mathcal{F}_e by unbounded operators. Moreover: they are not even a basis, while they are both complete in \mathcal{H} . The proofs of these claims do not differ much from those given in [18] and therefore will not be repeated here. We will comment further on the physical meaning of these results in the next subsection.

Similarly as in [18] we can prove that \mathcal{F}_{φ} and \mathcal{F}_{Ψ} are \mathcal{D} -quasi bases. In fact, repeating almost the same steps, we deduce that for instance, $\forall f, g \in \mathcal{D}$,

$$\langle f, g \rangle = \sum_{\underline{n}} \langle f, \varphi_{\underline{n}} \rangle \langle \Psi_{\underline{n}}, g \rangle, \quad (3.26)$$

so that the results listed at the end of section 2 hold true. In particular, let us introduce the operator $\Theta(\underline{\nu}, \underline{\mu}) := T(\underline{\mu}, \underline{\nu})$. It is possible to show that $\Theta(\underline{\nu}, \underline{\mu})$ is self-adjoint, invertible, and leaves \mathcal{D} invariant. Moreover, $\Theta(\underline{\nu}, \underline{\nu}) = 1$, and

$$\Theta(\underline{\nu}, \underline{\mu}) = N_{\Psi} \prod_{i=1}^2 e^{(\nu_i - \mu_i) \hat{a}_i^\dagger} e^{(\bar{\nu}_i - \bar{\mu}_i) \hat{a}_i}, \quad (3.27)$$

which implies that $\langle f, \Theta(\underline{\nu}, \underline{\mu})f \rangle > 0$ for all non zero vectors $f \in \mathcal{D}$. This is in agreement with the facts that (i) $\Psi_{\underline{n}} = \Theta(\underline{\nu}, \underline{\mu})\varphi_{\underline{n}}$, $\forall \underline{n}$; (ii) (a_j, b_j^\dagger) are Θ -conjugate: $a_j f = \Theta^{-1}(\underline{\nu}, \underline{\mu})b_j^\dagger \Theta(\underline{\nu}, \underline{\mu})f$, for all $f \in \mathcal{D}$. We conclude also that, again for all $f \in \mathcal{D}$,

$$N_i f = \Theta^{-1}(\underline{\nu}, \underline{\mu})N_i^\dagger \Theta(\underline{\nu}, \underline{\mu})f, \quad (3.28)$$

which is the intertwining relation responsible for the fact that \tilde{H} and \tilde{H}^\dagger have the same eigenvalues and related eigenvectors, see below.

¹This aspect is almost never stressed in the physical literature. Unbounded operators never exist *alone!* They exist in connection with some suitable dense subspace of \mathcal{H} , their domains.

3.2 Back to the Hamiltonian

Let us now return to our original problem, i.e. the deduction of the eigenvalues and the eigenvectors for \tilde{H} in (3.2) and \hat{H} in (1.1). As we have shown we may express them in terms of pseudo-bosonic number operators. From the above construction is clear that

$$\tilde{H}\varphi_{\underline{n}} = E_{\underline{n}}\varphi_{\underline{n}}, \quad (3.29)$$

with $E_{\underline{n}} \in \mathbb{R}$ given by (3.13). From our results in section 2 it also follows directly that the eigensystem of the adjoint $\tilde{H}^\dagger = \bar{\gamma}_1 N_1^\dagger + \bar{\gamma}_2 N_2^\dagger + \bar{\gamma}_0$ is computed to

$$\tilde{H}^\dagger \Psi_{\underline{n}} = \bar{E}_{\underline{n}} \Psi_{\underline{n}} = E_{\underline{n}} \Psi_{\underline{n}}. \quad (3.30)$$

The analysis in [18] showed that, as already deduced, two biorthogonal sets of eigenstates of a Hamiltonian and of its adjoint, need not to be automatically a Riesz basis, even when they are complete! This is exactly the case here: \mathcal{F}_φ and \mathcal{F}_Ψ are biorthogonal, complete, eigenstates of \tilde{H} and \tilde{H}^\dagger (\hat{H} and \hat{H}^\dagger), respectively, but neither \mathcal{F}_φ nor \mathcal{F}_Ψ are bases for \mathcal{H} . However, interestingly enough, they are \mathcal{D} -quasi bases, and this is reflected in the properties we have explicitly verified for our model.

4. Conclusions

We have investigated the properties of a non self-adjoint model on a noncommutative two dimensional space. The Hamiltonian \tilde{H} was set up in the standard fashion followed in the literature on \mathcal{PT} -symmetric quantum mechanics, by seeking an anti-linear symmetry, i.e. \mathcal{PT}_- in this case. From our explicit formulae we observe that \mathcal{PT}_- : $\varphi_{\underline{0}} \rightarrow \varphi_{\underline{0}}$, $\varphi_{\underline{n}} \rightarrow (-1)^{n_1} \varphi_{\underline{n}}$, $\Psi_{\underline{0}} \rightarrow \Psi_{\underline{0}}$, $\Psi_{\underline{n}} \rightarrow (-1)^{n_1} \Psi_{\underline{n}}$ such that by the standard arguments of Wigner [13] it follows that the eigenvalues of \tilde{H} have to be real. This is confirmed by our explicit computation. The symmetry for the Hamiltonian \hat{H} is not evident from the start, but as demonstrated the overall conclusions are the same as for \tilde{H} .

However, despite having well defined real physical spectrum, we established further that \tilde{H} can not be considered as a standard quantum mechanical model, since the corresponding biorthonormal system is not of Riesz type. As already discussed, in many places in the literature, see [19] for instance, it is incorrectly assumed that the eigenvectors of a not self-adjoint Hamiltonian H and H^\dagger automatically form a biorthogonal basis. In fact, this is a rather strong requirement which is quite difficult to find satisfied in concrete models existing in the literature, at least for infinite dimensional Hilbert spaces. We have shown that even for the simple example presented here this is not the case. This only leaves two of the following options: either this conclusion is wrong for the cases treated, as it would be for the model presented here, or at least some additional analysis is required to justify it. Thus our example supports the suggestion [14, 18] that many models, thought to be very interesting quantum mechanical systems, need to be revisited for further scrutiny.

It is easy to see from our formulae that these conclusions do not rely on the fact that the model is formulated on a noncommutative space and also hold in the limit to the commutative space when setting $\lim_{\theta \rightarrow 0} \Omega = 2\hbar$, $\lim_{\theta \rightarrow 0} M = m$, etc. In reverse, this

also means that the problem of not having automatically a biorthonormal basis can not be solved by formulating the model on a non-commutative space, which provides more freedom and often removes inconsistencies.

We end this section, and the paper, observing that, even with all the problems we have put in evidence along the paper, we may still make sense of the model presented here, simply because of the role of the quasi-bases as described above and in more detail in the quoted literature.

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References

- [1] C. M. Bender and S. Boettcher, Real Spectra in Non-Hermitian Hamiltonians Having PT Symmetry, *Phys. Rev. Lett.* **80**, 5243–5246 (1998).
- [2] A. Guo, G. J. Salamo, D. Duchesne, R. Morandotti, M. Volatier-Ravat, V. Aimez, G. A. Siviloglou and D.N. Christodoulides, Observation of PT -Symmetry Breaking in Complex Optical Potentials, *Phys. Rev. Lett.* **103**, 093902 (2009).
- [3] C. E. Ruter, K. G. Makris, R. El-Ganainy, D. N. Christodoulides, M. Segev, and D. Kip, Observation of parity-time symmetry in optics, *Nat. Phys.* **6**, 1515 (2010).
- [4] K. F. Zhao, M. Schaden and Z. Wu, Enhanced magnetic resonance signal of spin-polarized Rb atoms near surfaces of coated cells, *Phys. Rev. A*, **81**, 042903 (2010).
- [5] J. Schindler, A. Li, M. C. Zheng, F. M. Ellis and T. Kottos, Experimental study of active LRC circuits with PT symmetries, *Phys. Rev. A*, **84**, 040101, (2011).
- [6] C. M. Bender, Making Sense of Non-Hermitian Hamiltonians, *Rep. Progr. Phys.*, **70**, 947-1018 (2007).
- [7] A. Mostafazadeh, Pseudo-Hermitian representation of Quantum Mechanics, *Int. J. Geom. Methods Mod. Phys.* **7**, 1191-1306 (2010).
- [8] M. Znojil, Three-Hilbert-space formulation of Quantum Mechanics, *SIGMA* **5**, 001 (2009).
- [9] J. Dieudonné, Quasi-hermitian operators, *Proceedings of the International Symposium on Linear Spaces, Jerusalem 1960*, Pergamon, Oxford, 115–122 (1961).
- [10] F. G. Scholtz, H. B. Geyer, F.J.W. Hahne, Quasi-hermitian operators in quantum mechanics and the variational principle, *Ann. Phys.* **213**, 74–101 (1992).
- [11] M. Froissart, Covariant formalism of a field with indefinite metric *Il Nuovo Cimento* **14**, 197–204 1959.
- [12] A. Mostafazadeh, Pseudo-hermiticity versus pt symmetry. the necessary condition for the reality of the spectrum, *J. Math. Phys.* **43**, 205–214 2002.
- [13] E. Wigner, Normal form of antiunitary operators, *J. Math. Phys.* **1**, 409–413 1960.
- [14] D. Krejcirik and P. Siegl, On the metric operator for the imaginary cubic oscillator, *Phys. Rev. D*, **86**, 121702(R) (2012).

- [15] D. A. Trifonov, Pseudo-Boson Coherent and Fock States, in Differential Geometry, Complex Analysis and Mathematical Physics, eds. K. Sekigawa et al (Singapore: World Scientific 2009), pp. 241-250; arXiv:0902.3744[quant-ph].
- [16] F. Bagarello, More mathematics for pseudo-bosons, J. Math. Phys. **54**, 063512 (2013) (11pg).
- [17] F. Bagarello and M. Lattuca, \mathcal{D} pseudo-bosons in quantum models, Phys. Lett. A, in press.
- [18] F. Bagarello, From self to non self-adjoint harmonic oscillators: physical consequences and mathematical pitfalls, Phys. Rev. A, in press.
- [19] C. M. Bender and P. D. Mannheim, Exactly solvable PT -symmetric Hamiltonian having no Hermitian counterpart, Phys. Rev. D **78**, 025022 (2008); G. Dattoli, R. Mignani, A. Torre, Geometrical phase in the cyclic evolution of non-Hermitian systems, J. Phys. A: Math. Gen. **23** 5795-5806 (1990); F. H. M. Faisal, J. V. Moloney, Time-dependent theory of non-Hermitian Schrodinger equation: application to multiphoton-induced ionisation decay of atoms, J. Phys. B: At. Mol. Phys. **14**, 3603-3620 (1981)
- [20] F. Bagarello, Deformed canonical (anti-)commutation relations and non hermitian hamiltonians, in *Non-selfadjoint operators in quantum physics: Mathematical aspects*, F. Bagarello, J. P. Gazeau, F. Szafraniek and M. Znojil Eds, J. Wiley and Sons, in preparation
- [21] S. Dey, A. Fring, and L. Gouba, PT-symmetric noncommutative spaces with minimal volume uncertainty relations, J. Phys. **A45**, 385302 (2012).
- [22] J.-Q. Li, Y.-G. Miao and Z. Xue, Algebraic method for pseudo-Hermitian Hamiltonians, arXiv:1107.4972 [quant-ph]