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# A LIEB-THIRRING INEQUALITY FOR EXTENDED ANYONS

THÉOTIME GIRARDOT AND NICOLAS ROUGERIE

ABSTRACT. We derive a Pauli exclusion principle for extended fermion-based anyons of any positive radius and any non-trivial statistics parameter. That is, we consider 2D fermionic particles coupled to magnetic flux tubes of non-zero radius, and prove a Lieb-Thirring inequality for the corresponding many-body kinetic energy operator. The implied constant is independent of the radius of the flux tubes, and proportional to the statistics parameter.

## CONTENTS

1. Introduction	1
1.1. Anyons	2
1.2. Lieb-Thirring inequalities for ideal anyons	3
1.3. Model for extended anyons	4
1.4. Main theorem	5
1.5. Strategy of proof	5
2. Reduction to local estimates with finite $N$	7
2.1. Reduction to local Lieb-Thirring at finite $N$	7
2.2. Reduction to a local exclusion principle	9
3. Local exclusion principle for extended anyons	10
3.1. Large boxes	13
3.2. Medium Boxes	18
3.3. Small boxes	20
3.4. Conclusion of proofs	26
References	28

## 1. INTRODUCTION

The exclusion principle of quantum physics can be formulated in terms of Lieb-Thirring inequalities for the kinetic energy of fermionic particles. These inequalities stay true in various contexts, for instance, when the particles feel an external magnetic field. In this paper we establish a Lieb-Thirring inequality for extended anyons, modeled as fermions coupled to magnetic flux tubes of finite radius. The magnitude  $\alpha \in [0, 2]$  of the magnetic flux is interpreted as  $1+$  the statistics parameter (because our basic wave-functions are fermionic). Our motivation is two-fold:

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- (1) We bridge a gap in the current state of the research program initiated in [32]. Indeed, in [31] a Lieb-Thirring inequality is established for any  $\alpha \neq 1$ , at zero radius. Thus, ideal anyons of any statistics except the bosonic one satisfy a Pauli exclusion principle. On the other hand, in [20], bounds suggestive of a Pauli principle are proven for finite radii  $R$ , but under restrictive assumptions on  $\alpha$  and  $R$ .
- (2) We improve the main results [14, Theorems 1.1 and 1.4] of a previous paper of ours. Indeed, conditioned on the inequality we prove below, a semi-classical effective model for almost-fermionic extended anyons could be derived under relaxed assumptions in a mean-field type limit. This was mentioned in [14, Remark 1.2] and proved<sup>1</sup> in the first author's phd thesis [13, Chapter 15].

Before stating our inequality precisely we quickly recall basic facts about the two main concepts of the paper: extended anyons and Lieb-Thirring inequalities. In particular we give precisions on the vocabulary used above. We will not consider the case of non-abelian anyons where the exchange phase is replaced by a general unitary operator. See [28] and references therein.

**1.1. Anyons.** In a quantum mechanical system, dimensionality plays a fundamental role. In three or higher dimensions, the indistinguishability of the particles naturally leads to sorting them out into two types, bosons and fermions. These two types have different statistical behavior, leading to the commonly used terminology of bosonic (or fermionic) *statistics*. This dichotomy no longer holds in two dimensions. The richer topology of the configuration space indeed allows for more than two statistics. Consider a wave function  $\Psi : (\mathbb{R}^2)^N \rightarrow \mathbb{C}$ , it will formally behave as

$$\Psi(\mathbf{x}_1, \dots, \mathbf{x}_j, \dots, \mathbf{x}_k, \dots, \mathbf{x}_N) = e^{i\alpha_b\pi} \Psi(\mathbf{x}_1, \dots, \mathbf{x}_k, \dots, \mathbf{x}_j, \dots, \mathbf{x}_N) \quad (1.1)$$

where  $\alpha_b \in [-1, 1]$  is the statistics parameter of the anyons, counted from the bosonic end. The case  $\alpha_b = 0$  corresponds to bosons and  $\alpha_b = 1$  to fermions. The possibility of different statistics have been known since the 70's from different approaches [22, 35, 43, 15] and used to describe quasi-particles emerging in the fractional quantum Hall effect [18, 6, 17, 2, 6, 16, 30], rotating Bose gases [44, 7] and in quantum information [38]. There are two main ways to model anyons. We can treat  $\Psi$  either as a multi-valued function (a section of a complex line bundle) or, as a usual bosonic or fermionic function with a modified kinetic energy. We follow the latter approach, called the magnetic gauge picture.

We thus consider a fermionic wave function whose kinetic energy is modified through a singular change of gauge. Namely, we encode the behaviour of the wave function (1.1) under a particle exchange by setting

$$\Psi(\mathbf{x}_1, \dots, \mathbf{x}_N) = \prod_{j < k} e^{i\alpha_f \phi_{jk}} \Phi(\mathbf{x}_1, \dots, \mathbf{x}_N) \quad \text{where } \phi_{jk} = \arg \frac{\mathbf{x}_j - \mathbf{x}_k}{|\mathbf{x}_j - \mathbf{x}_k|}$$

with  $\Phi$  a fermionic wave function, antisymmetric under particle exchange and

$$\alpha_f = 1 + \alpha_b \quad (1.2)$$

the statistics parameter, now counted from the fermionic end. We have denoted  $\arg(\cdot)$  the angle of a planar vector with the horizontal axis. The case  $\alpha_f = 0$  describes a usual

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<sup>1</sup>The final improvement is actually from  $\eta < 1/4$  to  $\eta < 1/3$  in the notation of these references, slightly less than we had hoped for in [14, Remark 1.2].

fermionic system. Applying this transformation, the momentum operator for particle  $j$  changes as

$$-i\nabla_{\mathbf{x}_j} \rightarrow D_j := -i\nabla_{\mathbf{x}_j} + \alpha_f \mathbf{A}(\mathbf{x}_j) \quad (1.3)$$

with

$$\mathbf{A}(\mathbf{x}_j) := \sum_{k \neq j} \frac{(\mathbf{x}_j - \mathbf{x}_k)^\perp}{|\mathbf{x}_j - \mathbf{x}_k|^2} \quad (1.4)$$

where  $(x, y)^\perp = (-y, x)$ . Namely, we have the formal identity

$$\langle \Psi | (-i\nabla_{\mathbf{x}_j})^2 \Psi \rangle = \langle \Phi | D_{\mathbf{x}_j}^2 \Phi \rangle.$$

The above is a description of ideal anyons: the system behaves like ordinary particles attached to infinitely thin solenoids perpendicular to the plane. In other words we added a magnetic Aharonov-Bohm type interaction between the particles, which can formally be gauged away by changing the symmetry type of wave-functions. This corresponds to the particular case of anyons of radius  $R = 0$ . In this work we are interested in deriving Lieb-Thirring inequalities for anyons of radius  $R > 0$ , meaning that the Aharonov-Bohm flux tube will be smeared over a finite radius. Before we describe this in details we recall what is known for the above model at  $R = 0$ . We refer to [1, 32, 33, 8, 9] for more details on the definition of the model, in particular different possible self-adjoint extensions. In the sequel we always use the Friedrichs extension.

**1.2. Lieb-Thirring inequalities for ideal anyons.** The celebrated Lieb-Thirring inequalities are one of the different ways to quantify the Pauli exclusion between fermionic particles. We know, for instance [24, Theorem 4.3] that, in two dimensions, fermions exclude one another in the sense that for any anti-symmetric wave-function  $\Psi$

$$\sum_{j=1}^N \int_{\mathbb{R}^{2N}} |\nabla_{\mathbf{x}_j} \Psi|^2 d\mathbf{x} \geq C_2^K \int_{\mathbb{R}^2} \rho(\mathbf{x})^2 d\mathbf{x} \quad (1.5)$$

where

$$\rho(\mathbf{x}) := \sum_{j=1}^N \int_{\mathbb{R}^{d(N-1)}} |\Psi(\mathbf{x}_1, \dots, \mathbf{x}_j = \mathbf{x}, \dots, \mathbf{x}_N)|^2 \prod_{k \neq j} d\mathbf{x}_k$$

is the one-particle density. This inequality remains true with  $-i\nabla \rightarrow -i\nabla + \mathbf{A}$  for some suitable magnetic vector potential  $\mathbf{A}$ . It implies that the kinetic energy of a large number  $N$  of fermionic particles inside a bounded domain  $\Omega \subset \mathbb{R}^2$  must grow proportionally to  $N^2 \gg N$ : fermions do not like to be packed together and make it known through the large energetic cost of such a packing.

On the other hand, the best we can achieve with bosons is

$$\sum_{j=1}^N \int_{\mathbb{R}^{dN}} |\nabla_{\mathbf{x}_j} \Psi|^2 d\mathbf{x} \geq \frac{C_2}{N} \int_{\mathbb{R}^d} \rho(\mathbf{x})^2 d\mathbf{x} \quad (1.6)$$

which is actually the Sobolev inequality and becomes trivial as  $N \rightarrow \infty$ . Hence the kinetic energy of a large number  $N$  of bosonic particles inside a bounded domain  $\Omega \subset \mathbb{R}^2$  may well stay of order  $N$ .

As regards anyons, we intuitively think that the closer they are to fermions, the more they exclude one another. This directly leads to the idea of a Lieb-Thirring inequality for

anyons, proportional to  $\alpha$ . Results in this direction are available in [32, 33, 34]. Here we quote the more recent [31] where a bound is obtained for any  $\alpha_b \in [-1, 1]$ :

**Theorem 1.1 (Lieb-Thirring inequality for ideal anyons).**

Let  $D_{\mathbf{x}_j}$  be as in (1.3). There exists a constant  $C > 0$  such that for any  $\alpha_f \in [0, 2]$ ,  $N \geq 1$  and  $\Psi_N \in L^2_{\text{asym}}(\mathbb{R}^{2N})$

$$\sum_{j=1}^N \int_{\mathbb{R}^{dN}} |D_{\mathbf{x}_j} \Psi_N|^2 dx \geq C |1 - \alpha_f| \int_{\mathbb{R}^d} \rho_{\Psi_N}(\mathbf{x})^2 d\mathbf{x}$$

Beware that we use the fermionic convention that wave-functions are anti-symmetric. In [31] the result is stated with  $\Psi_N \in L^2_{\text{sym}}(\mathbb{R}^{2N})$  and  $1 - \alpha_f$  replaced by  $\alpha_b$  as per (1.2). This inequality shows that anyonic particles of any type but the bosonic one ( $\alpha_b = 0, \alpha_f = 1$ ) satisfy a Lieb-Thirring inequality and thus an exclusion principle.

The usual approach to Lieb-Thirring inequalities for fermions is to see them as dual to bounds on eigenvalue sums for Schrödinger operators, and use the Birman-Schwinger principle, see [24, Chapter 4] for review. This clearly does not apply in the anyonic case, because the problem is genuinely many-body. A new approach based on the local exclusion came up in the past ten years [11, 34, 32, 27, 31, 26, 21, 37, 36]. It consists in proving possibly  $N$ -dependent inequalities on finite subsets of  $\mathbb{R}^d$ . A clever covering of the space then allows to patch the inequalities together and obtain the correct  $N$ -dependence on the whole space. We will employ this technique here, applying it to the Hamiltonian for extended anyons that we describe next (see also [20, 29, 12, 14]).

**1.3. Model for exented anyons.** Consider the 2D Coulomb potential generated by a unit charge smeared over the disk of radius  $R$

$$w_R(\mathbf{x}) = (\log |\cdot| * \chi_R)(\mathbf{x}), \text{ with the convention } w_0 = \log |\cdot| \quad (1.7)$$

and  $\chi_R(x)$  a positive, regularizing function of unit mass

$$\chi_R(\mathbf{x}) := \frac{\mathbb{1}_{B(0,R)}(\mathbf{x})}{\pi R^2}. \quad (1.8)$$

Observe that

$$\nabla^\perp w_0(\mathbf{x}) = \frac{\mathbf{x}^\perp}{|\mathbf{x}|^2}, \text{ and } B_0(\mathbf{x}) = \nabla^\perp \nabla^\perp w_0 = \Delta w_0 = 2\pi\delta_0$$

so that we recover the magnetic field of the ideal anyon case (in a distributional sense) at  $R = 0$ . A natural regularisation of the ideal anyons potential vector (1.4) is

$$\mathbf{A}^R(\mathbf{x}_j) := \sum_{k \neq j} \nabla^\perp w_R(\mathbf{x}_j - \mathbf{x}_k) := \sum_{k \neq j} \frac{(\mathbf{x}_j - \mathbf{x}_k)^\perp}{|\mathbf{x}_j - \mathbf{x}_k|_R^2}. \quad (1.9)$$

where we have introduced the regularized distance

$$|\mathbf{x}|_R := \max\{|\mathbf{x}|, R\}.$$

The magnetic field felt by particle  $j$  is then

$$\text{curl} \mathbf{A}^R(\mathbf{x}_j) = 2\pi \sum_{k \neq j} \frac{\mathbb{1}_{B(\mathbf{x}_k, R)}(\mathbf{x}_j)}{\pi R^2},$$

i.e. it sees all the other particles as carrying a tube of flux of radius  $R$ . From now on we only use the fermionic representation and set

$$\alpha := \alpha_f = 1 + \alpha_b.$$

We always assume that

$$\alpha \in [0, 2].$$

The full kinetic energy operator is

$$T_\alpha^R := \sum_{j=1}^N \left( D_{\mathbf{x}_j}^R \right)^2 := \sum_{j=1}^N \left( -i\nabla_{\mathbf{x}_j} + \alpha \mathbf{A}^R(\mathbf{x}_j) \right)^2 \quad (1.10)$$

acting on the fermionic space  $L_{\text{asym}}^2(\mathbb{R}^{2N})$  as an unbounded operator. We denote  $\mathcal{D}_{\alpha,R}^N$  the domain of (1.10). When  $R > 0$ ,  $\mathbf{A}^R$  is a bounded perturbation of  $-i\nabla$ . The kinetic energy  $T_\alpha^R$  is essentially self-adjoint on its natural domain (see [39, Theorem X.17] and [3]). The bottom of its spectrum exists for any fixed  $R > 0$ .

Apart from providing an analytically useful regularisation of the model, the above Hamiltonian with smeared flux tubes is actually the relevant one for emergent anyons [30, 19] in the fractional quantum Hall effect. The size of the flux tubes is set by the magnetic length of the host system. Early considerations of the model are in [42, 41].

**1.4. Main theorem.** We define the kinetic energy of the  $N$ -particles system with wave function  $\Psi_N$  as

$$\mathcal{E}_\alpha^R[\Psi_N] = \sum_{j=1}^N \int_{\mathbb{R}^{2N}} \left| (-i\nabla_{\mathbf{x}_j} + \alpha \mathbf{A}^R(\mathbf{x}_j)) \Psi_N \right|^2 d\mathbf{x}_1 \dots d\mathbf{x}_N. \quad (1.11)$$

We also denote the one-body density of the system

$$\rho_{\Psi_N}(\mathbf{x}) := \sum_{j=1}^N \int_{\mathbb{R}^{d(N-1)}} |\Psi_N(\mathbf{x}_1, \dots, \mathbf{x}_j = \mathbf{x}, \dots, \mathbf{x}_N)|^2 \prod_{k \neq j} d\mathbf{x}_k$$

and state the main theorem of the paper.

**Theorem 1.2 (Lieb-Thirring inequality for extended anyons).**

*There exists a constant  $C^{\text{EA}}$  independent of  $R$ ,  $\alpha$  and  $N$  such that for any  $L^2$ -normalized fermionic  $N$ -particles wave function  $\Psi_N \in L_{\text{asym}}^2(\mathbb{R}^{2N})$*

$$\mathcal{E}_\alpha^R[\Psi_N] \geq C^{\text{EA}} |\alpha - 1| \int_{\mathbb{R}^2} \rho_{\Psi_N}^2(\mathbf{x}) d\mathbf{x}.$$

**1.5. Strategy of proof.** The rest of the paper is dedicated to the proof of Theorem 1.2. We follow the route opened in [32, 33, 34] where, roughly speaking, the authors divide the plane into squares  $Q$ , of various sizes  $|Q| = L^2$ , on which they prove a local Lieb-Thirring inequality before recombining. To do so, one needs

- (1) **A local exclusion principle** stating that, on a given square, the presence of two anyons<sup>2</sup> is sufficient to get a lower bound on the kinetic energy proportional to  $L^{-2}$ . For usual fermions this just means that the kinetic energy can have only one zero

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<sup>2</sup>Of course, the property “a square contains a certain number of particles” is a probabilistic statement.

mode per box (namely, the constant function), so that only one particle per box can have zero kinetic energy.

- (2) **A local uncertainty principle** derived from Poincaré-Sobolev inequalities. This does not use the statistics and is equally valid for bosons. Combining with the local exclusion yields a local Lieb-Thirring inequality on squares with more than two particles.
- (3) **A smart splitting algorithm** ensuring that the total energy of the squares with more than two particles is sufficient to compensate for squares on which we do not have enough particles to obtain the inequality.

In the sequel we essentially keep the same framework with the main following steps:

- (1) **A local exclusion principle for extended anyons** stating that if a box contains more than a fixed number  $\underline{N}$  of particles, its energy must be positive, proportional to  $L^{-2}$ . The proof technique we use rather depends on the ratio  $\gamma = R/L$ .
- (2) **A local uncertainty principle** derived from the diamagnetic and Sobolev inequalities and quite similar to the previously mentioned one.
- (3) **The Besicovitch covering theorem** allowing us to cover the plane with sets each containing sufficiently many particles to apply the local exclusion principle, while intersecting one another only a finite number of times.

The third idea was introduced in the recent [37], and allows to think purely locally, without having to look for compensations<sup>3</sup> between different spatial regions as in [32, 33, 34]. Hence our main task is to provide the local exclusion principle.

If  $\gamma = R/L \ll 1$  (large boxes), the fact that the anyons are extended intuitively does not play a very big role, and one should be able to adapt arguments from [31, 20] to obtain the exclusion. The difference is that in [31] the influence of particles outside of the box can be gauged away freely because the attached magnetic flux is purely local. We prove that for  $\gamma \ll 1$  this influence can be gauged away at a small, controlable cost.

If  $\gamma = R/L \sim 1$  (medium boxes) we can use the well-known (combine [23, Theorem 7.21] and [10, Lemma 1.4.1])

$$\langle \psi, (-i\nabla + A)^2 \psi \rangle \geq \frac{1}{2} \langle |\psi|, (-\Delta + \text{curl}A)|\psi| \rangle \quad (1.12)$$

to obtain bounds using a two-body model (the curl of (1.9) is a pair interaction). For  $\gamma \sim 1$  the Dyson lemma [25, 40] allows to use the kinetic energy to smear the two-body interaction over the whole box and get a non-trivial lower bound. This argument is worked out in [20], whose results we quote and adapt to our situation.

If  $\gamma = R/L \gg 1$  (small boxes), the box is completely covered by the magnetic flux attached to each particle inside it. We can then show that the problem becomes effectively one-body. Intuitively, the magnetic field does not harm the fact that the kinetic energy only has one zero-mode on the box. We prove a diamagnetic bound vindicating that there are only finitely many modes with energy less than  $L^{-2}$ , uniformly in the magnetic field. Since our wave-functions are fermionic, this implies the needed lower bound if sufficiently many particles are in the box.

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<sup>3</sup>We definitely could have used the splitting algorithms of [32, 26] in our proof instead of the Besicovitch theorem.

Strictly speaking (1.12) is not available on a box with Neumann boundary conditions. Hence, throughout the paper we apply it first on the whole space to obtain

$$\begin{aligned} \mathcal{E}_\alpha^R[\Psi_N] &\geq \int_{\mathbb{R}^{2N}} \frac{1}{2} \sum_{j=1}^N |(-i\nabla_{\mathbf{x}_j} + \alpha \mathbf{A}^R(\mathbf{x}_j)) \Psi_N|^2 d\mathbf{x}_1 \dots d\mathbf{x}_N \\ &\quad + \int_{\mathbb{R}^{2N}} \frac{1}{4} \left( \sum_{j=1}^N |\nabla_{\mathbf{x}_j} \Psi_N|^2 + 2\pi \sum_{k \neq j} \frac{\mathbb{1}_{B(\mathbf{x}_k, R)}(\mathbf{x}_j)}{\pi R^2} |\Psi_N|^2 \right) d\mathbf{x}_1 \dots d\mathbf{x}_N \\ &=: \int_{\mathbb{R}^{2N}} (e_1(\Psi_N) + e_2(\Psi_N)) d\mathbf{x}_1 \dots d\mathbf{x}_N \end{aligned} \quad (1.13)$$

and derive lower bounds to the right-hand side for the above.

In Section 2 we state our local exclusion bound and explain how to deduce Theorem 1.2 using local uncertainty and the Besicovitch theorem, as in [37]. The heart of the paper is then Section 3 where we prove the local exclusion estimate, distinguishing according to the size of the box. All in all, the logical structure of the argument is

$$\text{Section 3} \Rightarrow \text{Theorem 2.3} \Rightarrow \text{Theorem 2.1} \Rightarrow \text{Theorem 1.2.} \quad (1.14)$$

We dispose of the last two implications first, with essentially known methods, in order to focus on the key new estimates in Section 3.

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## 2. REDUCTION TO LOCAL ESTIMATES WITH FINITE $N$

As usual with Lieb-Thirring inequalities, the main point of Theorem 1.2 is the optimal dependence on  $N$ . In this section we reduce the proof to local estimates with unspecified  $N$  dependence via the Besicovitch covering theorem, following [37].

**2.1. Reduction to local Lieb-Thirring at finite  $N$ .** Let  $Q \subset \mathbb{R}^2$  a bounded domain (always taken to be a square in the sequel). We denote the local kinetic energy in  $Q$  by

$$\mathcal{E}_Q^R[\Psi_N] := \sum_{j=1}^N \int_{\mathbb{R}^{2N}} (e_1(\Psi_N) + e_2(\Psi_N)) \mathbb{1}_Q(\mathbf{x}_j) dx_N \quad (2.1)$$

with the energy densities  $e_1, e_2$  as defined in (1.13).

Our local Lieb-Thirring inequality, to be derived in Section 2.2 is as follows

**Theorem 2.1 (Lieb-Thirring at finite  $N$ ).**

*There exists three numbers  $N_<$ ,  $N_>$  and  $C^{\text{FN}}$  such that for any  $\Psi_N \in L_{\text{asym}}^2(\mathbb{R}^{2N})$  with*

$$N_< \leq \int_Q \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x} \leq N_> \quad (2.2)$$

and any square  $Q$  we have that

$$\mathcal{E}_Q^R[\Psi_N] \geq C^{\text{FN}} |\alpha - 1| \int_Q \rho_{\Psi_N}^2(\mathbf{x}) d\mathbf{x}.$$

We now explain how this implies Theorem 1.2. We recall the Besicovitch covering lemma [4, 5] which was used to prove Lieb-Thirring inequalities in [37] (using balls instead of squares).

**Lemma 2.2 (Besicovitch covering lemma).**

Let  $E$  be a bounded subset of  $\mathbb{R}^d$ . Let  $\mathcal{F}$  be a collection of hypercubes in  $\mathbb{R}^d$  with faces parallel to the coordinate planes such that every point  $\mathbf{x} \in E$  is the center of a cube from  $\mathcal{F}$ . Then there exist a sub-collection  $\mathcal{G} \subset \mathcal{F}$  such that

$$\mathbb{1}_E \leq \sum_{Q \in \mathcal{G}} \mathbb{1}_Q \leq b_d \mathbb{1}_E, \quad (2.3)$$

namely,  $E$  is covered by  $\bigcup_{Q \in \mathcal{G}} Q$  and every point in  $E$  belongs to at most  $b_d$  cubes from  $\mathcal{G}$ . The constant  $b_d$  only depends on the dimension  $d \geq 1$ .

*Proof of Theorem 1.2, last implication in (1.14).*

We start by considering  $\Psi_N$ 's which do not satisfy the assumption in Equation (2.2) of Theorem 2.1, i.e, with

$$N_{<} > \int_{\mathbb{R}^2} \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x}. \quad (2.4)$$

In this case, we can apply the diamagnetic [23, Theorem 7.21] and Sobolev inequalities

$$\mathcal{E}_\alpha^R[\Psi_N] \geq \sum_{j=1}^N \int_{\mathbb{R}^{2N}} |-\mathrm{i}\nabla_{\mathbf{x}_j} |\Psi_N||^2 d\mathbf{x}_N \geq \frac{1}{C_2} \frac{\int_{\mathbb{R}^2} \rho_{\Psi_N}^2(\mathbf{x}) d\mathbf{x}}{\int_{\mathbb{R}^2} \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x}} \geq \frac{1}{C_2 N_{<}} \int_{\mathbb{R}^2} \rho_{\Psi_N}^2(\mathbf{x}) d\mathbf{x}.$$

It remains to consider the case

$$N_{<} \leq \int_{\mathbb{R}^2} \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x}.$$

By a density argument we may assume that  $\Psi_N$  is smooth, with compact support in  $E^N$  with  $E \subset \mathbb{R}^2$  bounded. Then  $\rho_{\Psi_N}$  is continuous with a bounded support  $E \subset \mathbb{R}^2$ , and we can for every  $\mathbf{x} \in E$ , find a square  $Q_{\mathbf{x}} \subset \mathbb{R}^2$  centered at  $\mathbf{x}$  such that

$$\int_{Q_{\mathbf{x}}} \rho_{\Psi_N}(\mathbf{y}) d\mathbf{y} = \frac{N_{<} + N_{>}}{2} \quad (2.5)$$

where  $N_{>}$  is as in the statement of Theorem 2.1. We apply the Besicovitch covering lemma 2.2 to the collection of squares  $\mathcal{F} = \{Q_{\mathbf{x}}\}_{\mathbf{x} \in E}$  to obtain a sub-collection  $\mathcal{G} \subset \mathcal{F}$  such that

$$\mathbb{1}_E \leq \sum_{Q \in \mathcal{G}} \mathbb{1}_Q \leq b_2 \mathbb{1}_E. \quad (2.6)$$

The second inequality in the above implies that

$$\mathcal{E}_\alpha^R[\Psi_N] \geq \frac{1}{b_2} \sum_{Q \in \mathcal{G}} \mathcal{E}_Q^R[\Psi_N].$$

On each square  $Q \in \mathcal{G}$  we have (2.5) and may thus apply Theorem 2.1 to obtain

$$\mathcal{E}_\alpha^R[\Psi_N] \geq \frac{C^{\text{FN}}|\alpha-1|}{b_2} \sum_{Q \in \mathcal{G}} \int_Q \rho_{\Psi_N}^2(\mathbf{x}) d\mathbf{x} \geq \frac{C^{\text{FN}}|\alpha-1|}{b_2} \int_{\mathbb{R}^2} \rho_{\Psi_N}^2(\mathbf{x}) d\mathbf{x}$$

where we used the first inequality of (2.6) in the last step. This provides the desired estimate with

$$C^{\text{EA}} = \min \left\{ \frac{C^{\text{FN}}}{b_2}, \frac{1}{C_2 N_<} \right\}.$$

□

**2.2. Reduction to a local exclusion principle.** We can now state the Local Pauli exclusion theorem we use to establish Theorem 2.1. Its proof will be the content of Section 3.

**Theorem 2.3 (Local exclusion principle for extended anyons).**

*There exists three numbers  $N_<$ ,  $N_>$  and  $C^{\text{EA}}$  independent of  $\alpha$  and  $R$  such that for any square  $Q$  of side-length  $L$  and any  $\Psi_N \in L^2_{\text{asym}}(Q^N)$  satisfying*

$$N_< \leq \int_Q \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x} \leq N_> \quad (2.7)$$

*we have that*

$$\mathcal{E}_Q^R[\Psi_N] \geq C^{\text{LE}} \frac{|\alpha-1|}{|Q|} \int_Q \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x}. \quad (2.8)$$

To prove Theorem 2.1 we combine the above with a local uncertainty principle, i.e. essentially a Poincaré-Sobolev inequality. We use the version from [36, Lemma 3.4], which is convenient for our purpose.

**Lemma 2.4 (Local uncertainty).**

*Let  $\Psi_N \in H^1(\mathbb{R}^{2N})$  for arbitrary  $N \geq 1$  and let  $Q$  be a square in  $\mathbb{R}^2$ . Then*

$$\mathcal{E}_Q^R[\Psi_N] \geq \frac{1}{4} \sum_{j=1}^N \int_{\mathbb{R}^{2N}} |\nabla_{\mathbf{x}_j} |\Psi_N||^2 \mathbf{1}_Q(\mathbf{x}_j) \geq \frac{C_2 \int_Q \rho_{\Psi_N}^2(\mathbf{x}) d\mathbf{x}}{\int_Q \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x}} - \frac{1}{|Q|} \int_Q \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x} \quad (2.9)$$

*for a universal constant  $C_2$ .*

*Proof.* We apply [36, Lemma 3.4] to the wave function  $|\Psi_N|$ . □

*Proof of Theorem 2.1, second implication in (1.14).* We assume (2.7). Combining Inequalities (2.8) and (2.9) we obtain, for any  $\varepsilon \in [0, 1]$

$$\begin{aligned} (1 - \varepsilon + \varepsilon) \mathcal{E}_Q^R[\Psi_N] &\geq \varepsilon \frac{C_2 \int_Q \rho_{\Psi_N}^2(\mathbf{x}) d\mathbf{x}}{\int_Q \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x}} - \frac{\varepsilon}{|Q|} \int_Q \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x} \\ &\quad + (1 - \varepsilon) \frac{C^{\text{LE}}|\alpha-1|}{|Q|} \int_Q \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x} \\ &\geq \varepsilon \frac{C_2}{N_>} \int_Q \rho_{\Psi_N}^2(\mathbf{x}) d\mathbf{x} + \frac{1}{|Q|} (C^{\text{LE}}|\alpha-1|N_< - \varepsilon (N_> + C^{\text{LE}}|\alpha-1|N_>)). \end{aligned}$$

We choose (clearly this is smaller than 1)

$$\varepsilon = \frac{C^{\text{LE}}|\alpha - 1|N_{<}}{N_{>} + C^{\text{LE}}|\alpha - 1|N_{>}}$$

and the desired result follows with

$$C^{\text{FN}} = \frac{C_2}{N_{>}} \frac{C^{\text{LE}}N_{<}}{N_{>} + C^{\text{LE}}|\alpha - 1|N_{>}} \geq \frac{C_2}{N_{>}^2} \frac{C^{\text{LE}}N_{<}}{1 + 2C^{\text{LE}}}.$$

□

### 3. LOCAL EXCLUSION PRINCIPLE FOR EXTENDED ANYONS

There remains to deal with the heart of the matter, namely the proof of Theorem 2.3. This result is true without any assumption on  $\gamma = RL^{-1}$ , but we use different proofs for three particular ranges of  $\gamma$ , thus covering all  $\gamma \in \mathbb{R}_+$ . We introduce two constants  $c_1$  and  $c_2 > c_1$  to be fixed later on and work in three types of boxes, corresponding to the sketch of proof given at the end of Section 1

- (1) **Large boxes** where  $\gamma < c_1$ .
- (2) **Medium boxes** where  $c_1 \leq \gamma \leq c_2$ .
- (3) **Small boxes** where  $\gamma > c_2$ .

In each case we establish a lower bound on the energy of  $n$  particles localized in  $Q$ , uniformly with respect to a number  $m \geq 0$  of particles outside the box. We combine the three results in the end of the section to obtain the energy  $\mathcal{E}_Q$  of (2.1). The logic is that the first bound will be proven for  $c_1$  small enough, the third bound for  $c_2$  large enough, but the second bound is valid for all values of  $c_1, c_2$ , provided that  $c_1$  be bounded away from zero and that  $c_2$  remains bounded above.

The following notation is used throughout this section. We define the Neumann ground-state energy for  $n$  extended anyons on a domain  $Q \subset \mathbb{R}^2$  interacting with  $m$  anyons in the exterior  $Q^c = \mathbb{R}^2 \setminus Q$ . We denote  $Q_0 = [0, 1]^2$  and define

$$\mathbf{A}_j^R(X_n, Y_m) = \sum_{\substack{k=1 \\ k \neq j}}^n \nabla^\perp w_R(\mathbf{x}_j - \mathbf{x}_k) + \sum_{k=1}^m \nabla^\perp w_R(\mathbf{x}_j - \mathbf{y}_k) \quad (3.1)$$

the magnetic vector potential for  $n$  particles living in  $Q$  (with coordinates  $X_n = (\mathbf{x}_1, \dots, \mathbf{x}_n)$ ) interacting with  $m$  fixed anyons located at  $Y_m = (\mathbf{y}_1, \dots, \mathbf{y}_m)$ , all outside  $Q$ . In line with the definition of  $e_2$  in (1.13) let also

$$V_R(X_n, Y_m) = 2\pi \sum_{k \neq j} \frac{\mathbb{1}_{B(\mathbf{x}_k, R)}(\mathbf{x}_j)}{\pi R^2} + 2\pi \sum_{j=1}^n \sum_{k=1}^m \frac{\mathbb{1}_{B(\mathbf{y}_k, R)}(\mathbf{x}_j)}{\pi R^2} \quad (3.2)$$

For  $\Psi_n \in L^2(Q^n)$ , set

$$\begin{aligned} \mathcal{E}_n^R(Q, Y_m)[\Psi_n] &:= \frac{1}{2} \sum_{j=1}^n \int_{Q^n} |(-i\nabla_{\mathbf{x}_j} + \alpha \mathbf{A}_j^R(X_n, Y_m)) \Psi_n|^2 dX_n \\ &\quad + \frac{1}{4} \sum_{j=1}^n \int_{Q^n} |\nabla_{\mathbf{x}_j} \Psi_n|^2 dX_n + \frac{1}{4} \int_{Q^n} V_R(X_n, Y_m) |\Psi_n|^2 dX_n \\ &=: \int_{Q^n} e(\Psi_n, Y_m) dX_n \end{aligned} \quad (3.3)$$

The fermionic Neumann energy is then

$$E_n^R(Q, m) := \inf_{Y_m \in (Q^c)^m} \inf \left\{ \mathcal{E}_n^R(Q, Y_m)[\Psi_n], \|\Psi_n\|_{L^2_{\text{asym}}(Q^n)} = 1 \right\}. \quad (3.4)$$

We will drop the arguments  $Q$  or  $m$  of the previous energy when  $Q = Q_0 = [0, 1]^2$  or  $m = 0$ , namely

$$E_n^R(m) := E_n^R(Q_0, m) \quad \text{and} \quad E_n^R := E_n^R(Q_0, 0).$$

We also introduce the notation

$$D_{\mathbf{x}_j}^R(Y_m) = -i\nabla_{\mathbf{x}_j} + \alpha \mathbf{A}_j^R(X_n, Y_m).$$

The main goal of this section is to prove

**Proposition 3.1 (Exclusion principle on finite boxes).**

There exists constants  $\underline{N} > 0$  and  $C > 0$  such that, for any  $R > 0$ ,  $m \geq 0$  and  $N \geq \underline{N}$

$$E_N^R(Q, m) \geq \frac{CN|1 - \alpha|}{|Q|}. \quad (3.5)$$

Before distinguishing between different box-sizes as alluded to above, we start by general considerations that will reduce the proof to bounded particle numbers. We borrow and/or adapt several arguments from [31].

**Lemma 3.2 (Scaling property of the energy on a square).**

For any square  $Q$  such that  $|Q| = L^2$  we have the scaling property

$$E_n^R(Q, m) = |Q|^{-1} E_n^{R/L}(m) \quad (3.6)$$

*Proof.* We first translate the variables to work on  $[0, L]^2$ , with  $L$  the side-length of the original square. We proceed to the change of variables  $\mathbf{x} \rightarrow L\mathbf{x}$ . The first term in (3.3) becomes

$$\sum_{j=1}^n \int_{Q_0^n} \left| \left( \frac{-i\nabla_{\mathbf{x}_j}}{L} + \alpha \mathbf{A}_j^R(X_n, Y_m) \right) \Psi(L\mathbf{x}_1, \dots, L\mathbf{x}_n) \right|^2 |Q|^n d\mathbf{x}_1 \dots d\mathbf{x}_n. \quad (3.7)$$

Using the definition (1.9), the property  $|L\mathbf{x}|_R = L|\mathbf{x}|_{R/L}$  this becomes

$$\frac{1}{|Q|} \sum_{j=1}^n \int_{Q_0^n} \left| \left( -i\nabla_{\mathbf{x}_j} + \alpha \mathbf{A}_j^{R/L}(X_n, \frac{Y_m}{L}) \right) \Phi(\mathbf{x}_1, \dots, \mathbf{x}_n) \right|^2 \quad (3.8)$$

with  $\Phi(\mathbf{x}_1, \dots, \mathbf{x}_n) = L^n \Psi(L\mathbf{x}_1, \dots, L\mathbf{x}_n)$  so that  $\int_{Q_0^n} |\Phi|^2 = 1$ . The second part of (3.3) is dealt with similarly and we conclude by taking the infima.  $\square$

Next we have the equivalent of [31, Lemma 4.2]

**Lemma 3.3 (Superadditivity).**

Let  $Q := \cup_{q=1}^K Q_q$  with  $\{Q_q\}_{q=1}^K$  a collection of disjoint and simply connected subsets of  $Q$ . Let  $\vec{n} \in \mathbb{N}^K$  such that  $\sum_{q=1}^K n_q = N$ . We define the potential

$$W(X_n) := \sum_{\vec{n}} \sum_{q=1}^K E_{n_q}^R(Q_q, m + n - n_q) \mathbb{1}_{\vec{n}}(\mathbf{x}_1, \dots, \mathbf{x}_N)$$

where  $\mathbb{1}_{\vec{n}}$  denote the characteristic function of the subset of  $Q^n$  where exactly  $n_q$  of the points  $\{\mathbf{x}_1, \dots, \mathbf{x}_N\}$  are in  $Q_q$  for all  $1 \leq q \leq K$ .

We have that

$$E_n^R(Q, m) \geq \int_{Q^n} W |\Psi|^2 \quad (3.9)$$

and in particular

$$E_n^R(Q, m) \geq \min_{\vec{n}} \sum_{q=1}^K E_{n_q}^R(Q_q, m + n - n_q) \quad (3.10)$$

*Proof.* Let  $\mathbb{1}_{\vec{n}}$  denote the characteristic function of the subset of  $Q^n$  where exactly  $n_q$  of the points  $\{\mathbf{x}_1, \dots, \mathbf{x}_N\}$  are in  $Q_q$  for all  $1 \leq q \leq K$ . Then

$$\begin{aligned} \sum_{j=1}^n \int_{Q^n} e(\Psi, Y_m) dX_n &= \sum_{\{A_k\}} \int_{Q_1^{|A_1|}} \int_{Q_2^{|A_2|}} \cdots \int_{Q_K^{|A_K|}} \sum_{q=1}^K \sum_{j \in A_q} e(\Psi, Y_m) dX_A \\ &\geq \sum_{\{A_k\}} \sum_{q=1}^K E_{|A_q|}^R(Q_q, m + n - n_q) \int_{Q_1^{|A_1|}} \int_{Q_2^{|A_2|}} \cdots \int_{Q_K^{|A_K|}} |\Psi|^2 dX_A \end{aligned}$$

Here the sum over  $\{A_k\}$  runs over all partitions of the particles into the sets  $Q_q$ . We used the notation

$$dX_A = dX_{A_1} \cdots dX_{A_K} \quad \text{with} \quad dX_{A_l} = \prod_{j \in A_l} d\mathbf{x}_j.$$

□

Now we reduce to finite particle numbers as announced. Denote

$$E(n, Q) := \inf_{R>0, m \in \mathbb{N}} E_n^R(Q, m). \quad (3.11)$$

**Lemma 3.4 (Reduction to finite particle numbers).**

Let  $\bar{N} = 4^k$  for some integer  $k \geq 1$  and  $N \geq \bar{N}$ . Assume that  $E(n, Q) > 0$  for all  $4^{k-1} + 1 \leq n \leq N$ . There exists a constant  $C_k > 0$ , independent of  $N$ , such that

$$E(N, Q) \geq C_k N \min \left( E(4^{k-1} + 1, Q), \dots, E(4^k, Q) \right). \quad (3.12)$$

The assumption  $E(n, Q) > 0$  for all  $4^{k-1} + 1 \leq n \leq N$  is only temporary. Its validity will follow from the considerations in the next subsections.

*Proof.* This follows exactly the proof of [31, Lemma 4.8]. Without loss (in view of Lemma 3.2) we work on the unit square and drop  $Q$  from the notation. Splitting it into four equally large squares and using Lemmas 3.3 and 3.2 we have

$$E(N) \geq 4 \min_{\vec{n}} \sum_{q=1}^4 E(n_q).$$

It is here important that we have taken the infimum over  $R$  and  $m$  in (3.11). At least one of the squares must contain  $n/4$  particles, hence, dropping the other terms, we obtain

$$E(N) \geq 4 \min_{j=\lfloor \frac{N}{4} \rfloor + 1, \dots, N} E(j). \quad (3.13)$$

We treat the case  $N = 4^l$  for  $l$  integer for simplicity (the generalization is straightforward, as in [31, Lemma 4.8]). Denoting

$$e_n := \min \{E(j), 4^n + 1 \leq j \leq 4^{n+1}\}$$

we obtain from (3.13) that

$$e_n \geq 4 \min(e_n, e_{n-1})$$

and, since our assumption implies  $e_n > 0$ ,

$$e_n \geq 4e_{n-1}.$$

Iterating a finite number of times we deduce

$$E(n) \geq 4^{l-k} \min \left( E(4^{k-1}) + 1, \dots, E(4^k) \right) \geq C_k n \min \left( E(4^{k-1}) + 1, \dots, E(4^k) \right)$$

where  $C_k$  depends only on  $k$ , thus on  $\bar{N}$ . □

**3.1. Large boxes.** Our exclusion principle on large boxes reads as follows:

**Proposition 3.5 (Exclusion principle on large boxes).**

*There exists a small  $0 < c_1 < \frac{1}{12}$ , a natural number  $\bar{N} > 0$  and a  $C_n > 0$  independent of  $R$ ,  $m$ ,  $\alpha$  or  $Q$  such that, for any  $\gamma = R/L \leq c_1$  and  $2 \leq n \leq \bar{N}$*

$$E_n^R(Q, m) \geq \frac{C|1 - \alpha|}{|Q|}. \quad (3.14)$$

The proof is inspired from [31]. It occupies the rest of the subsection and requires several preparatory lemmas.

**Lemma 3.6 (Upper-bound for the two-particles energy on the unit square).**

*The two-particles energy on the unit square can be bounded by a constant  $C$  which does not depend on  $R$ ,  $m$  or  $\alpha$ .*

$$E_2^R(m) \leq C$$

*Proof.* We can construct a fermionic trial state on the unit cube  $Q_0$  by setting

$$\psi_2(\mathbf{x}_1, \mathbf{x}_2) := (x_1^1 - x_2^1) + (x_1^2 - x_2^2)$$

with the notation  $\mathbf{x}_j = (x_j^1, x_j^2)$ . We then have

$$\begin{aligned} \|\psi_2\|_{L^2(Q_0^2)} E_2^R(m) &\leq \mathcal{E}_2^R[\psi_2] = 2 \int_{Q_0^2} \left| -i\nabla_{\mathbf{x}_1} \psi_2 + \alpha \frac{(\mathbf{x}_1 - \mathbf{x}_2)^\perp}{|\mathbf{x}_1 - \mathbf{x}_2|_R^2} \psi_2 \right|^2 d\mathbf{x}_1 d\mathbf{x}_2 \\ &= 2 \int_{Q_0^2} \left( |\nabla_{\mathbf{x}_1} \psi_2|^2 + \alpha^2 \frac{|\psi_2|^2}{|\mathbf{x}_1 - \mathbf{x}_2|^2} \right) d\mathbf{x}_1 d\mathbf{x}_2 \\ &\leq 8 \end{aligned}$$

where we expanded the square using that  $\psi_2$  is real and applied both that  $|\alpha| \leq 1$  and that  $|\mathbf{x}_1 - \mathbf{x}_2|_R \geq |\mathbf{x}_1 - \mathbf{x}_2|$ .  $\square$

A key step is to reduce the lower bound to the energy with exactly two particles in a square. The basic idea is simple: if the ground state for  $n$  particles is nearly constant, it assigns non-zero probability to the event ‘‘a subsquare contains exactly two particles’’. If the ground state is not nearly constant, it must come with some positive kinetic energy.

**Lemma 3.7 (A priori bounds in terms of  $E_2^R$ ).**

There exist two constants  $C_1$  and  $C_2$  independent of  $R$ ,  $\alpha$ ,  $n$  or  $m$  such that

$$E_n^R(m) \geq \frac{C_n}{C_1 + C_2 C_n} E_2^{2R}(m + n - 2) \quad (3.15)$$

where  $C_n = \binom{n}{2} \left(\frac{3}{4}\right)^{n-2}$ .

*Proof.* We follow the same route as in [31, Lemma 4.3] but with fermionic wave functions becoming bosonic wave functions via the diamagnetic inequality.

We start from  $E_n^R(m)$  and divide  $Q_0$  of side-length 1 in four equally large squares of side-length  $1/2$

$$Q_0 := Q_1 \sqcup Q_2 \sqcup Q_3 \sqcup Q_4$$

We apply Lemma 3.3 to obtain

$$\sum_{j=1}^N \int_{Q_0^n} e(\Psi, Y_m) \geq \int_{Q_0^n} W |\Psi|^2$$

and we use the scaling property (3.6) to get

$$W \geq W_2 = 4E_2^{2R}(m + n - 2) \sum_{\vec{n}} \sum_{q=1}^4 (n_q = 2) \mathbb{1}_{\vec{n}}$$

with the notation  $(P) = 1$  when the statement  $P$  is true and  $(P) = 0$  otherwise. We can compute

$$\int_{Q_0^n} W_2 = E_2^R(m) \binom{n}{2} \left(\frac{3}{4}\right)^{n-2} = C_n E_2^R(m)$$

by counting the probability that exactly two particles are in a given square. We now want to estimate

$$\int_{Q_0^n} W_2 |\Psi|.$$

To this aim we will use a little bit of the kinetic energy and the inequality (3.9) to obtain

$$T_\alpha^{R,m} = \kappa T_\alpha^{R,m} + (1 - \kappa) T_\alpha^{R,m} \geq \kappa T_\alpha^{R,m} + (1 - \kappa) W_2$$

for any  $\kappa \in [0, 1]$ . The diamagnetic inequality leads to

$$\begin{aligned} E_n^R(m) &= \inf_{\Psi \in L_{\text{asym}}^2(Q_0^n), \|\psi\|=1} \langle \Psi, T_\alpha^{R,M} \Psi \rangle \\ &\geq \inf_{\Psi \in L_{\text{asym}}^2(Q_0^n), \|\psi\|=1} \langle |\Psi|, (-\kappa \Delta_{Q_0^n} + (1 - \kappa) W_2) |\Psi| \rangle \\ &\geq \inf_{\Psi \in L_{\text{sym}}^2(Q_0^n), \|\psi\|=1} \text{spec}(-\kappa \Delta_{Q_0^n} + (1 - \kappa) W_2) \end{aligned} \quad (3.16)$$

where  $\Delta_{Q^N}$  is the neumann Laplacian on  $Q^N$ . Note that  $E_n^R(Q_0, m)$  was defined on fermionic wave functions whereas now the operator of (3.16) acts on bosons. We consider the orthogonal projection

$$P := |u_0\rangle\langle u_0|$$

onto the normalised ground state of  $\Delta_{Q^N}$ , i.e. the constant function  $u_0 \equiv 1$ , and the orthogonal complement

$$P^\perp := \mathbb{1} - P$$

for which we have

$$-\Delta_{Q_0^n} \geq \pi^2 P^\perp.$$

We use a Cauchy-Schwarz inequality to get that

$$W_2 = (P + P^\perp) W_2 (P + P^\perp) \geq (1 - \varepsilon) P W_2 P + (1 - \varepsilon^{-1}) P^\perp W_2 P^\perp$$

for arbitrary  $\varepsilon \in [0, 1]$ .

Now we have

$$P W_2 P = P \int_{Q_0^n} W_2$$

and

$$P^\perp W_2 P^\perp \leq P^\perp \|W_2\|_\infty \leq 16 E_2^{2R}(m) P^\perp.$$

The combination of the previous estimates leads to

$$\begin{aligned} -\kappa \Delta_{Q_0^n} + (1 - \kappa) W_2 &\geq (\kappa \pi^2 - (1 - \kappa) (\varepsilon^{-1} - 1)) 16 E_2^{2R}(m) P^\perp \\ &\quad + (1 - \kappa) (1 - \varepsilon) C_n E_2^{2R}(m) P. \end{aligned}$$

We choose  $\kappa$  to make the prefactors in front of the projections equal

$$\kappa = \frac{E_2^{2R}(m) [C_n (1 - \varepsilon) + 16 (\varepsilon^{-1} - 1)]}{\pi^2 + E_2^{2R}(m) [C_n (1 - \varepsilon) + 16 (\varepsilon^{-1} - 1)]}$$

We then obtain the bound

$$E_n^R(m) \geq \frac{\pi^2 C_n (1 - \varepsilon) E_2^{2R}(m)}{\pi^2 + E_2^{2R}(m) [C_n (1 - \varepsilon) + 16 (\varepsilon^{-1} - 1)]}$$

where the choice  $\varepsilon = \frac{1}{2}$  combined with the upper bound on  $E_2^R(m)$  of Lemma 3.6 gives the result.  $\square$

The above lemma provides a lower bound in terms of the energy with exactly two particles inside the box, but possibly with the extra influence of many fixed particles outside the box. For  $R = 0$ , the influence of the latter can be gauged away as in [31]. In our case, the magnetic flux-tube they carry might overlap the box. We next check they still can be gauged away, at a controllable cost. Indeed, since  $R \ll L$  the flux-tube of outside particles touches the particles in the box only when they are close to the boundary.

**Lemma 3.8 (Gauging away particles outside the box).**

*There exists a small  $\gamma_0 < \frac{1}{4}$  such that for any  $\gamma \leq \gamma_0$  we can find a constant  $C_1$  independent of  $R$ ,  $m$ ,  $\alpha$  or  $Q$  and a constant  $C_{\gamma_0}$  only depending on  $\gamma_0$  such that*

$$\text{either } E_2^R(Q, m) \geq C_1 E_2^R(Q) \text{ or } E_2^R(Q, m) \geq \frac{C_{\gamma_0}}{|Q|}. \quad (3.17)$$

*Proof.* Here we drop the second term in (3.3) to consider only the magnetic kinetic energy. We remove from  $Q$  a very thin corridor such that the particles (of small radius) inside the restricted domain cannot interact with the outside. The proof will show that if the density is small in the corridor, we can neglect it. If not then the kinetic energy has to be large enough for the statement to hold.

We assume that  $L \geq 4R$  and define  $S = [2R, L - 2R]^2$ . On  $S^{\times 2}$ , define the change of gauge

$$\tilde{\Psi}_2 = \prod_{k=1}^m e^{i\alpha(\phi_{1k} + \phi_{2k})} \Psi_2 \quad \text{with } \phi_{jk} := \arg \frac{\mathbf{x}_j - \mathbf{y}_k}{|\mathbf{x}_j - \mathbf{y}_k|} \quad (3.18)$$

where  $\mathbf{y}_1, \dots, \mathbf{y}_m$  are the coordinates of particles outside the box. Note that

$$\nabla_{\mathbf{x}_j} \phi_{jk} = \frac{(\mathbf{x}_j - \mathbf{y}_k)^\perp}{|\mathbf{x}_j - \mathbf{y}_k|^2} = \frac{(\mathbf{x}_j - \mathbf{y}_k)^\perp}{|\mathbf{x}_j - \mathbf{y}_k|_R^2}$$

when  $|\mathbf{x}_j - \mathbf{y}_k| \geq R$ , and that  $\nabla_{\mathbf{x}_j} \phi_{jk}$  is regular and curl-free in that range. We thus have, using also the fermionic symmetry,

$$\begin{aligned} \mathcal{E}_2^R(Q, Y_m)[\Psi_2] &\geq \int_{Q^2} \left| \left( -i\nabla_1 + \alpha \frac{(\mathbf{x}_1 - \mathbf{x}_2)^\perp}{|\mathbf{x}_1 - \mathbf{x}_2|_R^2} + \alpha \sum_{k=1}^m \frac{(\mathbf{x}_1 - \mathbf{y}_k)^\perp}{|\mathbf{x}_1 - \mathbf{y}_k|_R^2} \right) \Psi_2 \right|^2 d\mathbf{x}_1 d\mathbf{x}_2 \\ &\geq \int_{Q^2} \left| \left( -i\nabla_1 + \frac{(\mathbf{x}_1 - \mathbf{x}_2)^\perp}{|\mathbf{x}_1 - \mathbf{x}_2|_R^2} \right) \tilde{\Psi}_2 \right|^2 \mathbb{1}_S(\mathbf{x}_1) d\mathbf{x}_1 d\mathbf{x}_2 \\ &\geq \frac{1}{2} E_2^R(Q) \int_{Q^2} |\Psi_2|^2 \mathbb{1}_S(\mathbf{x}_1) d\mathbf{x}_1 d\mathbf{x}_2 \\ &\geq \frac{1}{2} E_2^R(Q) \left( \int_{Q^2} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 - \int_{Q \times \Omega} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 \right) \end{aligned} \quad (3.19)$$

where we denote  $\Omega = Q \setminus S$  and use that  $|\tilde{\Psi}_2| = |\Psi_2|$ . Note that

$$|\Omega| = 8LR - 16R^2 \leq 8LR \quad (3.20)$$

We then observe that if for a small number  $c < 1$

$$\int_{Q \times \Omega} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 \leq c \int_{Q^2} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2$$

the first inequality of (3.17) results from (3.19) with  $C_1 = \frac{1}{2}(1 - c)$ .

On the other hand, when

$$\int_{Q \times \Omega} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 \geq c \int_{Q^2} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 \quad (3.21)$$

we introduce  $P$ , the orthogonal projector on the constant function and  $P^\perp$  such that

$$P + P^\perp = \mathbf{1}_{L^2(Q^4)}.$$

Then

$$\begin{aligned} \int_{Q \times \Omega} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 &\leq 2 \frac{|Q \times \Omega|}{|Q \times Q|} \int_{Q^2} |P|\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 + 2 \int_{Q \times \Omega} |P^\perp|\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 \\ &\leq 16 \frac{L^3 R}{L^4} \int_{Q^2} |P|\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 + 2|Q \times \Omega|^{\frac{1}{2}} \left( \int_{Q^2} |P^\perp|\Psi_2|^4 d\mathbf{x}_1 d\mathbf{x}_2 \right)^{\frac{1}{2}} \end{aligned}$$

where we used the Cauchy-Schwarz inequality and (3.20). We now apply a Poincaré-Sobolev inequality [23, Theorem 8.11] to  $P^\perp|\Psi_2| = |\Psi_2| - P|\Psi_2|$

$$\left( \int_{Q^2} |P^\perp|\Psi_2|^4 d\mathbf{x}_1 d\mathbf{x}_2 \right)^{\frac{1}{2}} \leq C_{\text{ps}} \int_{Q^2} |\nabla|\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2$$

which provides

$$\begin{aligned} C_{\text{ps}} \int_{Q^2} |\nabla|\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 &\geq \frac{1}{8L\sqrt{LR}} \int_{Q \times \Omega} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 - \frac{R}{2L^2\sqrt{LR}} \int_{Q^2} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 \quad (3.22) \\ &\geq \frac{1}{L^2} \int_{Q^2} |\Psi_2|^2 d\mathbf{x}_1 d\mathbf{x}_2 \left( c(4\sqrt{\gamma})^{-1} - \sqrt{\gamma} \right) \end{aligned}$$

using (3.20) again and Assumption (3.21) for the first term in the right hand side of (3.22). The diamagnetic inequality [23, Theorem 7.21] allows to obtain the second bound of (3.17) provided that we chose a  $\gamma$  small enough.  $\square$

We are now reduced to a bound on the energy of two particles inside the box, with no particles outside. This can be estimated using previous results of [20]. Essentially, with two isolated particles the effect of the flux tube between them may never be spoiled<sup>4</sup> by that of other particles, and it provides a non-zero energy via a Hardy-type inequality.

We need to introduce the function

$$\begin{aligned} g^2 : \mathbb{R}_+ \times [0, 1] &\rightarrow \mathbb{R} \\ \nu, \gamma &\mapsto g^2(\nu, \gamma) \end{aligned}$$

which is the smallest positive solution  $\lambda$  associated with the Bessel equation

$$-u'' - \frac{u'}{r} + \nu^2 \frac{u}{r^2} = \lambda u \quad (3.23)$$

on the interval  $[\gamma, 1]$  with Neumann boundary conditions while  $g(\nu, \gamma) = \nu$  for  $\gamma \geq 1$ . The function  $g$  has the property that

$$g(\nu, \gamma) \xrightarrow{\gamma \rightarrow 0} j'_\nu \geq \sqrt{2\nu} \quad (3.24)$$

<sup>4</sup>Recall that an integer flux can be gauged away.

where  $j'_\nu$  denotes the first positive zero of the derivative of the Bessel function  $J_\nu$ .

**Lemma 3.9 (Bound on  $E_2^R$ ).**

There exists a  $\gamma_1 < 1/12$  such that for all  $\gamma \leq \gamma_1$  there is a constant  $C_{\gamma_1}$  only depending on  $\gamma_1$  such that

$$E_2^\gamma \geq C_{\gamma_1} |\alpha - 1|. \quad (3.25)$$

*Proof.* We apply [20, Lemma 5.3]. Note that this result is stated for anyons based on a bosonic wave function while in our case we used a fermionic one (1.1). However, the proof given in [20] may be applied to this case. Indeed, the basic building block is [20, Lemma 3.1]. We use it in the antipodal-antisymmetric case rather than the antipodal-symmetric case, and otherwise follow the rest of the proof mutatis mutandis. The result is

$$E_2^R \geq \frac{\pi}{48} g^2 (c\alpha_2, 12\gamma) (1 - 12\gamma)_+^3$$

where

$$\alpha_N := \min_{p \in \{0, 1, \dots, N-2\}} \min_{q \in \mathbb{Z}} |(2p+1)(1-\alpha) - 2q|.$$

We use that  $\alpha_2 = |1 - \alpha|$  for  $\alpha \in [0, 2]$ . Combining with the limit property (3.24) we can pick a  $\gamma_1$  such that

$$E_2^R \geq C_{\gamma_1} |1 - \alpha|.$$

□

We may now conclude the

*Proof of Proposition 3.5.* We start by applying Lemma 3.7 with the scaling property (3.6) to the energy

$$E_n^R(Q, m) \geq \frac{C_n}{|Q|} E_2^{2\gamma} (m + n - 2).$$

We use Lemma 3.8 to bound the two-particles energy appearing in the above. We then restrict  $2\gamma \leq \gamma_0$  and obtain that

$$E_n^R(Q, m) \geq \frac{C_n}{|Q|} \times \left\{ C_1 E_2^{2\gamma} \text{ or } C_{\gamma_0} \right\}.$$

If we consider the  $C_{\gamma_0}$  case there is nothing more to prove and the final constant is  $C_{\gamma_0} C_n$ . In the other case we use Lemma 3.9 which provides a new restriction  $2\gamma \leq \gamma_1$  under which the bound (3.25) leads to

$$E_n^R(Q, m) \geq \frac{C_n}{|Q|} C_1 C_{\gamma_1} |1 - \alpha|.$$

This concludes the proof, upon redefining  $C_n$  and taking  $c_1 = \frac{1}{2} \min\{\gamma_0, \gamma_1\}$ . The final constant is not uniform in  $n$ , whence our assumption that  $n \leq \overline{N}$ . □

**3.2. Medium Boxes.** We now deal with medium boxes where  $c_1 \leq \gamma \leq c_2$  with  $c_1 < \sqrt{2} \leq c_2$ . In this case [20, Lemma 5.1] essentially provides the desired bound.

**Proposition 3.10 (Exclusion principle on medium boxes).**

Let  $c_1 \leq \gamma = R/L \leq c_2$  and  $Q$  be a square of side-length  $L$ . There is a constant  $C(c_1, c_2)$  depending only on  $c_1, c_2$  such that

$$E_n^R(m, Q) \geq C(c_1, c_2) \frac{|1 - \alpha| (n - 1)_+}{L^2}.$$

*Proof.* We simplify notation by working on the unit square. We introduce the function

$$K_\alpha = \sqrt{2|\alpha|} \frac{I_0(\sqrt{2|\alpha|})}{I_1(\sqrt{2|\alpha|})}$$

where  $I_\nu$  is the modified Bessel function of order  $\nu$ . One can show that

$$CI_0(2) \geq K_\alpha \geq 2$$

when  $\alpha \in [0, 2]$ . Indeed, the second bound comes from [20, Lemma 5.1] while the first one follows from the fact that  $I_0$  is an increasing function and that for  $x \in [0, 2]$ ,

$$I_1(x) := \sum_{m=0}^{\infty} \frac{1}{m! \Gamma(m+2)} \left(\frac{x}{2}\right)^{2m+1} \geq \frac{x}{2\Gamma(2)}.$$

We use [20, Lemma 5.1], which provides two bounds depending on the range of  $\gamma$ , stated for bosonic based anyons. The proof starts from (1.12), so that lower bounds are obtained in terms of the modulus of the wave-function, which is always bosonic. One can thus follow the argument mutatis mutandis in our case.

The first bound of the lemma holds when  $\gamma < \sqrt{2}$  and is obtained by the application of Dyson's lemma on the kinetic energy added to the magnetic interaction energy obtained via (1.12) (see [20, Lemma 1.1] for details). It states that

$$E_n^R(m) \geq \frac{|\alpha| \min\left\{(1 - \gamma^2/2)^{-1}, K_\alpha/2\right\}}{K_\alpha + 2|\alpha|(-\ln(\gamma/\sqrt{2}))} (n-1)_+.$$

Under the additional assumption that  $c_1 \leq \gamma$  the divergence of the logarithm is under control and the bound reduces to

$$E_n^R(m) \geq C \frac{|\alpha|}{I_0(\sqrt{2}) + (-\log(c_1/\sqrt{2}))} (n-1)_+ \geq C_1 |\alpha| (n-1)_+.$$

The second bound is valid for any  $\gamma \geq \sqrt{2}$  and is obtained using the magnetic interaction of [20, Lemma 1.1] where the indicator function equals 1 on the whole box  $Q$ :

$$E_n^R(m) \geq 2|\alpha|\gamma^{-2}n(n-1)_+ \geq \frac{2|\alpha|}{c_2^2} (n-1)_+.$$

These bounds are conclusive when  $|\alpha| \geq c > 0$ , choosing

$$C(c_1, c_2) = \min\{C_1, 2/c_2^2\}.$$

There remains to deal with the case where  $|\alpha|$  is allowed to become small, were we have to use that our basic wave-functions are fermionic, a constraint we dropped in the argument above (indeed, we have used only the second term in (3.3)).

Consider the function

$$h(\alpha) := \inf_{R, n \geq 2, m \in \mathbb{N}} \frac{E_n^R(m)}{(n-1)_+}$$

where the infimum on  $R$  is taken over a finite sub-interval of  $(0, \infty)$ , as in the statement of the result. The function  $h$  is an infimum of continuous functions of  $\alpha$  and is thus itself continuous. On the other hand, the above considerations imply that

$$h(\alpha) \geq C|\alpha|.$$

Moreover, at  $\alpha = 0$ ,  $h(\alpha)$  is bounded below in terms of the grounds-state problem for free fermions in a box, hence  $h(0) \geq c > 0$ . It follows by continuity that  $h(\alpha) \geq c' > 0$  for any  $\alpha$ , which concludes the proof.  $\square$

**3.3. Small boxes.** We now work under the assumption that  $\gamma \geq c_2 > \sqrt{2}$ . We shall actually assume more, namely that the particle number is uniformly bounded from above, and that  $\gamma$  is sufficiently large compared to that upper bound. We use only the first term in (3.3) to obtain:

**Proposition 3.11 (Exclusion principle on small boxes).**

Assume that  $\gamma \geq c_2 \geq \sqrt{2}$ . There exists a  $\underline{N}$  such that, if  $\underline{N} < n \leq N$  and  $R/L \geq C\sqrt{N}$  for a sufficiently large constant  $C$ , then

$$E_n^R(Q, m) \geq C \frac{n}{|Q|} \quad (3.26)$$

independently of  $R$  and  $m$ .

The proof, given at the end of the subsection, will consist of the combination of several lemmas that we discuss first.

The magnetic field created by the particles inside of the box is constant for  $\gamma > \sqrt{2}$ , covering the whole box, and the total flux in the box is small when  $\gamma$  is large. One can then reduce to a fermionic one-body problem with an external magnetic field.

**Lemma 3.12 (Reduction to a one-particle operator).**

Assume  $R > \sqrt{2}L$ , let  $Y_m = (\mathbf{y}_1, \dots, \mathbf{y}_m) \in (Q^c)^{\times m}$ . Define the vector potential

$$\mathbf{A}(\mathbf{x}) := (n-1) \frac{\mathbf{x}^\perp}{R^2} + \sum_{j=1}^m \nabla^\perp w_R(\mathbf{x} - \mathbf{y}_j) \quad (3.27)$$

associated to the magnetic field

$$\text{curl } \mathbf{A}(\mathbf{x}) = \frac{2}{R^2} \left( (n-1) + \sum_{k=1}^m \mathbb{1}_{|\mathbf{x} - \mathbf{y}_j| \leq R} \right). \quad (3.28)$$

Define the (Neumann realization of the) associated one-particle magnetic Schrödinger operator

$$h_{\mathbf{A}} := (-i\nabla_{\mathbf{x}} + \alpha \mathbf{A}(\mathbf{x}))^2 = \sum_{k=1}^{\infty} \lambda_k |u_k\rangle \langle u_k| \quad (3.29)$$

together with its spectral decomposition (with eigenvalues labeled in increasing order and  $L^2$ -normalized eigenfunctions). We have that, for any  $\Psi_n \in L_{\text{asym}}^2(Q^n)$  and any  $k \in \mathbb{N}$

$$\mathcal{E}_n^R(Q, Y_m) [\Psi_n] \geq \frac{1}{2} n \lambda_k \left( 1 - \frac{k-1}{n} - C \frac{knL^2}{R^2} \right). \quad (3.30)$$

*Proof.* We use only the first term in (1.13) here. By fermionic symmetry

$$2\mathcal{E}_n^R(Q, Y_m) [\Psi_n] \geq n \int_{Q^n} \left| \left( -i\nabla_{\mathbf{x}_1} + \alpha \sum_{k=2}^n \nabla^\perp w_R(\mathbf{x}_1 - \mathbf{x}_k) + \alpha \sum_{j=1}^m \nabla^\perp w_R(\mathbf{x}_1 - \mathbf{y}_j) \right) \Psi_n \right|^2 d\mathbf{x}_1 \dots d\mathbf{x}_n.$$

We observe that, under our assumption on  $\gamma = R/L$  we have

$$\nabla^\perp w_R(\mathbf{x}_1 - \mathbf{x}_k) = \frac{(\mathbf{x}_1 - \mathbf{x}_k)^\perp}{R^2}$$

for  $\mathbf{x}_1, \mathbf{x}_k \in Q$ . Define then

$$\varphi(\mathbf{x}, \mathbf{y}) := -\frac{\mathbf{x} \cdot \mathbf{y}^\perp}{R^2} \quad \text{and} \quad \Phi_n := \prod_{k=2}^n e^{-i\alpha\varphi(\mathbf{x}_1, \mathbf{x}_k)} \Psi_n$$

Since

$$\nabla_{\mathbf{x}} \varphi(\mathbf{x}, \mathbf{y}) = -\frac{\mathbf{y}^\perp}{R^2}$$

we find that

$$\begin{aligned} \int_{Q^n} \left| \left( -i\nabla_{\mathbf{x}_1} + \alpha \sum_{k=2}^n \nabla^\perp w_R(\mathbf{x}_1 - \mathbf{x}_k) + \alpha \sum_{j=1}^m \nabla^\perp w_R(\mathbf{x}_1 - \mathbf{y}_j) \right) \Psi_n \right|^2 d\mathbf{x}_1 \dots d\mathbf{x}_n = \\ \int_{Q^n} \left| \left( -i\nabla_{\mathbf{x}_1} + \alpha(n-1) \frac{\mathbf{x}_1^\perp}{R^2} + \alpha \sum_{j=1}^m \nabla^\perp w_R(\mathbf{x}_1 - \mathbf{y}_j) \right) \Phi_n \right|^2 d\mathbf{x}_1 \dots d\mathbf{x}_n = \\ \int_{Q^n} |(-i\nabla_{\mathbf{x}_1} + \alpha \mathbf{A}(\mathbf{x}_1)) \Phi_n|^2 d\mathbf{x}_1 \dots d\mathbf{x}_n. \end{aligned}$$

Inserting now the spectral decomposition of  $h_{\mathbf{A}}$

$$\begin{aligned} \int_{Q^n} |(-i\nabla_{\mathbf{x}_1} + \alpha \mathbf{A}(\mathbf{x}_1)) \Phi_n|^2 d\mathbf{x}_1 \dots d\mathbf{x}_n = \sum_{j=1}^{\infty} \lambda_j \int_{Q^{n-1}} \left| \int_Q \overline{u_j(\mathbf{x}_1)} \Phi_n(\mathbf{x}_1, \dots, \mathbf{x}_n) d\mathbf{x}_1 \right|^2 d\mathbf{x}_2 \dots d\mathbf{x}_n \\ \geq \lambda_k \left( 1 - \sum_{j=1}^{k-1} \int_{Q^{n-1}} \left| \int_Q \overline{u_j(\mathbf{x}_1)} \Phi_n(\mathbf{x}_1, \dots, \mathbf{x}_n) d\mathbf{x}_1 \right|^2 d\mathbf{x}_2 \dots d\mathbf{x}_n \right) \end{aligned}$$

where we picked an energy level  $\lambda_k$  and dropped the contribution of all levels with  $j < k$ .

Next we split

$$\begin{aligned} \int_{Q^{n-1}} \left| \int_Q \overline{u_j(\mathbf{x}_1)} \Phi_n(\mathbf{x}_1, \dots, \mathbf{x}_n) d\mathbf{x}_1 \right|^2 d\mathbf{x}_2 \dots d\mathbf{x}_n \\ = \int_{Q^n} \overline{u_j(\mathbf{x}_1)} u_j(\mathbf{y}_1) \Phi_n(\mathbf{x}_1, \dots, \mathbf{x}_n) \overline{\Phi_n(\mathbf{y}_1, \dots, \mathbf{x}_n)} d\mathbf{y}_1 d\mathbf{x}_1 d\mathbf{x}_2 \dots d\mathbf{x}_n = \text{I} + \text{II} \end{aligned}$$

with

$$\begin{aligned} \text{I} &= \int_{Q^n} \overline{u_j(\mathbf{x}_1)} u_j(\mathbf{y}_1) \Psi_n(\mathbf{x}_1, \dots, \mathbf{x}_n) \overline{\Psi_n(\mathbf{y}_1, \dots, \mathbf{x}_n)} d\mathbf{y}_1 d\mathbf{x}_1 d\mathbf{x}_2 \dots d\mathbf{x}_n \\ \text{II} &= \int_{Q^n} \left( \prod_{k=2}^n e^{i\alpha\varphi(\mathbf{x}_1, \mathbf{x}_k) - i\alpha\varphi(\mathbf{y}_1, \mathbf{x}_k)} - 1 \right) \overline{u_j(\mathbf{x}_1)} u_j(\mathbf{y}_1) \Psi_n(\mathbf{x}_1, \dots, \mathbf{x}_n) \overline{\Psi_n(\mathbf{y}_1, \dots, \mathbf{x}_n)} d\mathbf{y}_1 d\mathbf{x}_1 d\mathbf{x}_2 \dots d\mathbf{x}_n. \end{aligned}$$

Now, the first term I is just  $n^{-1}$  times the occupation number of the mode  $u_j$  in the *fermionic* wave-function  $\Psi_n$ . Hence

$$|\text{I}| \leq \frac{1}{n}.$$

On the other hand, since all variables sit within the box  $Q$  of side-length  $L$  we clearly have

$$|\varphi(\mathbf{x}_1, \mathbf{x}_k) - \varphi(\mathbf{y}_1, \mathbf{x}_k)| \leq C \frac{L^2}{R^2}$$

and hence

$$\left| 1 - \prod_{k=2}^n e^{i\alpha\varphi(\mathbf{x}_1, \mathbf{x}_k) - i\alpha\varphi(\mathbf{y}_1, \mathbf{x}_k)} \right| \leq C|\alpha| \frac{nL^2}{R^2}.$$

Thus, with a Cauchy-Schwartz inequality

$$\begin{aligned} |\text{II}| &\leq C \frac{nL^2}{R^2} \int_{Q^n} |u_j(\mathbf{x}_1)| |u_j(\mathbf{y}_1)| |\Psi_n(\mathbf{x}_1, \dots, \mathbf{x}_n)| |\Psi_n(\mathbf{y}_1, \dots, \mathbf{x}_n)| d\mathbf{y}_1 d\mathbf{x}_1 d\mathbf{x}_2 \dots d\mathbf{x}_n \\ &\leq C \frac{nL^2}{R^2} \int_{Q^n} |u_j(\mathbf{x}_1)|^2 |\Psi_n(\mathbf{y}_1, \dots, \mathbf{x}_n)|^2 d\mathbf{y}_1 d\mathbf{x}_1 d\mathbf{x}_2 \dots d\mathbf{x}_n = C \frac{nL^2}{R^2}. \end{aligned}$$

The desired final result is obtained by collecting the previous inequalities.  $\square$

Now we may look for lower bounds to a fermionic one-body problem with an external magnetic field. The latter can be quite general, because of the influence of particles sitting outside the box. Diamagnetic considerations however imply estimates independent of this field:

**Lemma 3.13 (Diamagnetic bound).**

*Consider the Neumann realization of the magnetic Laplacian*

$$H_{\mathbf{A}} := (-i\nabla + \mathbf{A})^2$$

*on  $H_{\mathbf{A}}^1(Q_0)$ , where  $\mathbf{A}$  and  $\text{curl}\mathbf{A}$  are bounded functions. Define the number*

$$N(\Lambda, \mathbf{A}) := \text{the number of eigenvalues of } H_{\mathbf{A}} \text{ less than or equal to } \Lambda.$$

*There exists a function  $f : \mathbb{R} \rightarrow \mathbb{R}$  independent of  $\mathbf{A}$  such that*

$$N(\Lambda, \mathbf{A}) \leq f(\Lambda).$$

We now turn to the proof of Lemma 3.13, first reducing to the case where  $\vec{\nu} \cdot \mathbf{A} \equiv 0$  on the boundary.

**Lemma 3.14 (Reduction to tangential vector potentials).**

*The general case of Lemma 3.13 is implied by the particular case where the normal component of  $\mathbf{A}$  vanishes on the boundary,  $\vec{\nu} \cdot \mathbf{A} \equiv 0$  on  $\partial Q$ .*

*Proof.* Let  $\phi$  be the unique solution to the Dirichlet problem

$$\begin{cases} \Delta\phi = B = \text{curl}(\mathbf{A}) \text{ in } Q_0 \\ \phi = 0 \text{ on } \partial Q_0. \end{cases}$$

By definition

$$\text{curl}\mathbf{A} = \text{curl}\nabla^\perp\phi$$

and hence there exists some  $\varphi$  such that

$$\mathbf{A} = \nabla^\perp \phi + \nabla \varphi \text{ in } Q_0.$$

Thus we may change gauge

$$e^{i\varphi} H_{\mathbf{A}} e^{-i\varphi} = H_{\nabla^\perp \phi}$$

where the multiplication operator  $e^{i\varphi}$  is unitary. Hence  $H_{\mathbf{A}}$  and  $H_{\nabla^\perp \phi}$  have the same spectrum. Since  $\phi$  is constant on the boundary, clearly  $\vec{\nu} \cdot \nabla^\perp \phi \equiv 0$  there.  $\square$

From now on we thus assume that

$$\begin{cases} \operatorname{curl} \mathbf{A} &= B \\ \mathbf{A} \cdot \vec{\nu} &= 0 \text{ on } \partial Q_0 \end{cases}$$

where  $\vec{\nu}$  is the normal vector of  $\partial Q_0$ . We denote  $\{\lambda_j(\mathbf{A})\}_{j=1}^\infty$  the eigenvalues of  $H_{\mathbf{A}}$ . The associated eigenfunctions  $\Psi_j$  are solutions of

$$\begin{cases} (-i\nabla + \mathbf{A})^2 \Psi_j &= \lambda_j \Psi_j \text{ on } Q_0 \\ \vec{\nu} \cdot (-i\nabla + \mathbf{A}) \Psi_j &= (-i\nabla \cdot \vec{\nu}) \Psi_j = 0 \text{ on } \partial Q_0. \end{cases}$$

For any real number  $e > 0$  we define the magnetic Neumann Green function  $G_{\mathbf{y}}^{\mathbf{A},e}$  to be the solution of

$$\begin{cases} (H_{\mathbf{A}} + e) G_{\mathbf{y}}^{\mathbf{A},e}(\mathbf{x}) &= \delta_{\mathbf{y}}(\mathbf{x}) \text{ on } Q_0 \\ (-i\nabla \cdot \vec{\nu}) G_{\mathbf{y}}^{\mathbf{A},e}(\mathbf{x}) &= 0 \text{ on } \partial Q_0. \end{cases}$$

in the weak sense that

$$\int_Q G_{\mathbf{y}}^{\mathbf{A},e}(\mathbf{x}) (-i\nabla + \mathbf{A})^2 \phi(\mathbf{x}) \, d\mathbf{x} + e \int_Q G_{\mathbf{y}}^{\mathbf{A},e}(\mathbf{x}) \phi(\mathbf{x}) \, d\mathbf{x} = \phi(\mathbf{y})$$

for all  $\phi \in C^\infty(Q)$ .

This way, the function

$$u_g(\mathbf{x}) = \int_Q G_{\mathbf{y}}^{\mathbf{A},e}(\mathbf{x}) g(\mathbf{y}) \, d\mathbf{y}$$

$g \in C^\infty(Q)$  is the unique solution of

$$(H_{\mathbf{A}} + e) u_g(\mathbf{x}) = g(\mathbf{x})$$

in the above weak sense (cf the Lax-Milgram theorem). In other words  $G_{\mathbf{y}}^{\mathbf{A},e}(\mathbf{x})$  is the integral kernel of  $(H_{\mathbf{A}} + e)^{-1}$ .

Our aim is now to show that the Green function with magnetic field is always smaller than the one without magnetic field. To this end we use Kato's inequality as in [24, Section 4.4]. We need a version thereof valid in the case of Neumann boundary conditions:

**Lemma 3.15 (Kato's Inequality with Neumann boundary conditions).**

Let  $u \in H^1(Q)$  such that  $(-i\nabla \cdot \vec{\nu}) u = 0$  and  $H_{\mathbf{A}} u$  be in  $L^1_{\text{loc}}(Q)$ . Let  $\operatorname{sgn}(u) = \bar{u}/|u|$  if  $u(\mathbf{x}) \neq 0$  and 0 otherwise. Then

$$-\Delta |u| \leq \operatorname{Re} [\operatorname{sgn}(u) H_{\mathbf{A}} u]$$

in the weak sense that the inequality holds when integrated against any positive  $\phi \in C^\infty(Q)$  with  $\vec{\nu} \cdot \nabla \phi \equiv 0$  on the boundary, i.e.

$$-\int_Q |u| \Delta \phi \leq \int_Q \operatorname{Re} [\operatorname{sgn}(u) H_{\mathbf{A}} u] \phi$$

*Proof.* We may follow the proof of [39, Theorem- X.33]. Essentially, if  $u$  is smooth and non-vanishing, the result holds pointwise. Where  $u$  vanishes,  $|u|$  is minimum, so that the left-hand side of the inequality is non-positive whereas the right-hand side is null by definition. Hence the inequality holds pointwise for smooth functions. Using a standard regularisation we deduce that the inequality holds in the weak sense by performing integrations by parts. The boundary terms vanish because we have  $\vec{\nu} \cdot \nabla \phi \equiv 0$  and  $\vec{\nu} \cdot \nabla |u| \equiv 0$  on the boundary.

A more complete proof is obtained as in [39, Theorems X.27 and X.33] by working with the regularized absolute value

$$u_\varepsilon(\mathbf{x}) = \sqrt{|u(\mathbf{x})|^2 + \varepsilon^2}.$$

The additional ingredient is the observation that  $\vec{\nu} \cdot \nabla u \equiv 0$  on the boundary implies that also  $\vec{\nu} \cdot \nabla u_\varepsilon \equiv 0$  on the boundary, so that integration by parts do not produce boundary terms. □

We prove another intermediary lemma.

**Lemma 3.16 (Positivity of the Laplace Green function).**

*The operator*

$$\begin{aligned} (-\Delta + e)^{-1} : L^2(Q) &\rightarrow H^1(Q) \\ f &\mapsto u \end{aligned}$$

*defined by*

$$\begin{cases} -\Delta u + eu = f & \text{on } Q \\ \vec{\nu} \cdot \nabla u = 0 & \text{on } \partial Q. \end{cases}$$

*preserves positivity, i.e.,  $(-\Delta + e)^{-1} f \geq 0$  if  $f \geq 0$ . Hence the Green function  $G^{0,e}$  with zero magnetic field is non-negative.*

*Proof.* The function  $u$  being the solution of  $-\Delta u + eu = f$  we know by the Lax-Milgram theorem that  $u$  is the unique minimizer on  $H^1(Q)$  of the energy

$$\mathcal{E}[u] = \frac{1}{2} \left( \int |\nabla u|^2 + e|u|^2 \right) - \int fu$$

but since  $f \geq 0$  we clearly have  $\mathcal{E}[u] \geq \mathcal{E}[|u|]$  and thus  $u = |u| \geq 0$  by uniqueness of the solution. □

We have all the ingredient to relate Green functions with and without magnetic field:

**Lemma 3.17 (Diamagnetic inequality for Green functions).**

*Let  $G_{\mathbf{x}}^{\mathbf{A},e}$  and  $G_{\mathbf{x}}^e = G_{\mathbf{x}}^{0,e}$  be defined as in (3.31). For almost every  $\mathbf{z} \in Q$*

$$|G_{\mathbf{x}}^{\mathbf{A},e}(\mathbf{z})| \leq G_{\mathbf{x}}^e(\mathbf{z})$$

*Proof.* We take two positive functions  $f, h \in C^\infty(Q)$  and define

$$u_h(\mathbf{x}) = \int_Q G_{\mathbf{y}}^{\mathbf{A},e}(\mathbf{x}) h(\mathbf{y}) d\mathbf{y} \quad \text{and} \quad u_f(\mathbf{x}) = \int_Q G_{\mathbf{y}}^e(\mathbf{x}) f(\mathbf{y}) d\mathbf{y}.$$

We know by Lemma 3.16 that  $u_f$  is positive. We apply Lemma 3.15 to  $u_h$  to obtain

$$(-\Delta + e)|u_h|(\mathbf{x}) \leq h(\mathbf{x}).$$

We now multiply by  $u_f$  and integrate to get that

$$\begin{aligned} \int_{Q^2} G_{\mathbf{y}}^e(\mathbf{x}) f(\mathbf{y}) (-\Delta_{\mathbf{x}} + e) |u_h(\mathbf{x})| \, d\mathbf{x}d\mathbf{y} &\leq \int_{Q^2} G_{\mathbf{y}}^e(\mathbf{x}) f(\mathbf{y}) h(\mathbf{x}) \, d\mathbf{x}d\mathbf{y} \\ \int_Q f(\mathbf{y}) |u_h(\mathbf{y})| \, d\mathbf{y} &\leq \int_{Q^2} G_{\mathbf{y}}^e(\mathbf{x}) f(\mathbf{y}) h(\mathbf{x}) \, d\mathbf{x}d\mathbf{y} \end{aligned}$$

where we used that  $(-\Delta_{\mathbf{x}} + e) G_{\mathbf{y}}^e(\mathbf{x}) = \delta_{\mathbf{y}}(\mathbf{x})$ . We then obtain (recalling that  $f \geq 0$ ),

$$\begin{aligned} \left| \int_{Q^2} G_{\mathbf{x}}^{\mathbf{A},e}(\mathbf{y}) f(\mathbf{y}) g(\mathbf{x}) \, d\mathbf{x}d\mathbf{y} \right| &\leq \int_{Q^2} G_{\mathbf{y}}^e(\mathbf{x}) f(\mathbf{y}) h(\mathbf{x}) \, d\mathbf{x}d\mathbf{y} \\ |G_{\mathbf{x}}^{\mathbf{A},e}(\mathbf{z})| &\leq G_{\mathbf{x}}^e(\mathbf{z}) \end{aligned}$$

by taking  $g \rightarrow \delta_{\mathbf{x}}$  and  $f \rightarrow \delta_{\mathbf{y}}$ . □

Now we may conclude the

*Proof of Lemma 3.13.* We define the Birman-Schwinger operator

$$K_e := \sqrt{\Lambda} (H^{\mathbf{A}} + e)^{-1} \sqrt{\Lambda}$$

If  $H^{\mathbf{A}}\psi = \lambda\psi$  we have

$$(H^{\mathbf{A}} + e)\psi = \Lambda\psi$$

with  $e = \Lambda - \lambda > 0$ . Hence  $\sqrt{\Lambda}\psi$  is the eigenfunction of  $K_e$  associated with the eigenvalue 1. We denote

$B_e :=$  the number of eigenvalues of  $K_e$  larger than or equal to 1.

The Birman-Schwinger principle (see [24, Equation 4.3.5]) states that

$$N(\Lambda, \mathbf{A}) = B_e. \tag{3.31}$$

We bound  $B_e$  in the following way, for any  $m \geq 1$

$$\begin{aligned} B_e &\leq \text{Tr} [K_e^m] \\ &= \Lambda^m \text{Tr} \left[ (H^{\mathbf{A}} + e)^{-m} \right] \\ &= \int_{Q^{2m}} G_e^{\mathbf{A}}(\mathbf{x}, \mathbf{y}_1) G_e^{\mathbf{A}}(\mathbf{y}_1, \mathbf{y}_2) \dots G_e^{\mathbf{A}}(\mathbf{y}_{m-1}, \mathbf{x}) \, d\mathbf{x}d\mathbf{y}_1d\mathbf{y}_2\dots d\mathbf{y}_{m-1} \end{aligned}$$

we take the absolute value and use Lemma (3.17) to get

$$\begin{aligned} B_e &\leq \Lambda^m \int_{Q^{2m}} G_e^0(\mathbf{x}, \mathbf{y}_1) G_e^0(\mathbf{y}_1, \mathbf{y}_2) \dots G_e^0(\mathbf{y}_{m-1}, \mathbf{x}) \, d\mathbf{x}d\mathbf{y}_1d\mathbf{y}_2\dots d\mathbf{y}_{m-1} \\ &= \Lambda^m \text{Tr} (-\Delta_{\mathbf{x}} + e)^{-m} \\ &= \Lambda^m \sum_{j=1}^{\infty} (\lambda_j(0) + \Lambda - \lambda)^{-m} \\ &=: f(\Lambda) \end{aligned} \tag{3.32}$$

where  $f$  is finite for  $m \geq 2$  and independent of  $\mathbf{A}$ . □

Finally we may give the

*Proof of Proposition 3.11.* Scaling length units by a factor  $L$ , we apply Lemma 3.13 to the magnetic Schrödinger operator (3.29). We thus find that the number of energy levels below a threshold  $\Lambda L^{-2}$  is bounded above by a function of  $\Lambda$  only. Pick some value of  $\Lambda$ , say  $\Lambda = 2$  for reference and denote  $\underline{N}$  the number of eigenvalues of  $h_A$  below  $\Lambda L^{-2}$ . Using (3.30) (with a proper choice of  $\lambda_k$  in relation with  $2L^{-2}$ ) and inserting the consequences of Lemma 3.13 we just mentioned, we find

$$\mathcal{E}_n^R(Q, Y_m) [\Psi_n] \geq CL^{-2}n (1 - (\underline{N} - 1)n^{-1} - C\underline{N}nL^2R^{-2}).$$

Under our stated assumptions on  $n$  and  $\gamma$ , the quantity inside the parenthesis is bounded below by a positive universal constant, and this concludes the proof.  $\square$

**3.4. Conclusion of proofs.** First note that Proposition 3.1 follows immediately from a combination of Lemma 3.4, Propositions 3.11-3.10-3.5:

*Proof of Proposition 3.1.* Let  $\underline{N}$  be as in Proposition 3.11 and  $\overline{N} = 4^k$  be such that  $\underline{N} \leq 4^{k-1} + 1$ . Using (3.11) and applying Lemma 3.4 we find

$$E_N^R(Q, m) \geq E(N, Q) \geq CN \min \left( E(4^{k-1} + 1, Q), \dots, E(4^k, Q) \right).$$

Note that the assumption  $E(n, Q) > 0$  for all  $4^{k-1} + 1 \leq n \leq N$  made in Lemma 3.4 is valid as shown by a combination of Propositions 3.11-3.10-3.5 and Lemma 3.7.

We are thus reduced to a uniform strictly positive lower bound on  $|Q|E(n)$  for particle numbers  $n$  with  $\underline{N} \leq n \leq \overline{N}$ . We obtain this using bounds on  $E(n, Q)$  provided by Propositions 3.5-3.10 or 3.11 depending on the value of  $\gamma = R/L$ . If  $\gamma \geq C(\overline{N})^{1/2}$  for  $C$  sufficiently large, we may use Proposition 3.11. We use (3.14) if  $\gamma < c_1$ , the constant in the statement of Proposition 3.5. Finally we use Proposition 3.10 in the remaining range of  $\gamma$ . In all cases we find

$$E(n, Q) = \inf_{R, m} E_n^R(Q, m) \geq \frac{C|1 - \alpha|}{|Q|}$$

when  $\underline{N} \leq n \leq \overline{N}$ , independently of  $R$ , and this yields the result.  $\square$

We next use Proposition 3.1 to conclude the

*Proof of Theorem 2.3.* We work under the assumption that  $\Psi_N$  has a density satisfying

$$1 + \underline{N} \leq N_{<} \leq \int_Q \rho_{\Psi_N}(\mathbf{x}) d\mathbf{x} \leq N_{>} \quad (3.33)$$

where  $\underline{N}$  is as in the statement of Proposition 3.1. We split the quantity to estimate according to how many particles are in the square  $Q$

$$\mathcal{E}_Q^R[\Psi_N] \geq \sum_{n=0}^N E_n^R(Q, m) p_n(\Psi_N, Q) \quad (3.34)$$

with  $p_n(\Psi_N, Q)$  the  $n$ -particle probability distribution induced from  $\Psi_N$

$$p_n(\Psi_N, Q) = \sum_{A \subseteq \{1, \dots, N\}, |A|=n} \int_{\overline{Q}^{N-n}} \int_{Q^n} |\Psi_N|^2 \prod_{k \in A} d\mathbf{x}_k \prod_{l \notin A} d\mathbf{x}_l$$

satisfying

$$\sum_{n=0}^N p_n(\Psi_N, Q) = 1 \quad \text{and} \quad \sum_{n=0}^N np_n(\Psi_N, Q) = \int_Q \rho_{\Psi_N}. \quad (3.35)$$

Then, using Proposition 3.1 we obtain

$$\mathcal{E}_Q^R[\Psi_N] \geq \frac{C|\alpha-1|}{|Q|} \sum_{n=\underline{N}}^N np_n(\Psi_N, Q).$$

On the other hand, using (3.33) and (3.35) we have

$$\begin{aligned} \sum_{n \geq \underline{N}} np_n(\Psi_N, Q) &\geq \underline{N} + 1 - \sum_{n < \underline{N}} np_n(\Psi_N, Q) \\ &\geq \underline{N} + 1 - (\underline{N} - 1) \sum_{n \geq 0} p_n(\Psi_N, Q) \\ &\geq 2 \end{aligned}$$

and hence

$$\mathcal{E}_Q^R[\Psi_N] \geq \frac{C|\alpha-1|}{|Q|N_{>}} \int_Q \rho_{\Psi_N}(\mathbf{x}) \, d\mathbf{x},$$

concluding the proof. □

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