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Rigidity of the Laughlin liquid

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We consider general N -particle wave functions that have the form of a product of the Laughlin state with filling factor $1/\ell$ and an analytic function of the N variables. This is the most general form of a wave function that can arise through a perturbation of the Laughlin state by external potentials or impurities, while staying in the lowest Landau level and maintaining the strong correlations of the original state. We show that the perturbation can only shift or lower the 1-particle density but nowhere increase it above a maximum value. Consequences of this bound for the response of the Laughlin state to external fields are discussed.

In theoretical studies of the fractional quantum Hall effect (FQHE) [9, 10, 12, 35] Laughlin's wave function(s) [14–16] play a fundamental role. There is such a function for every positive integer ℓ and it can be written, in units where the magnetic length is $1/\sqrt{2}$, as

$$\Psi_{\text{Lau}} = c_{\text{Lau}} \prod_{i < j} (z_i - z_j)^\ell e^{-\sum_{i=1}^N |z_i|^2/2} \quad (1)$$

where the $z_i \in \mathbb{C}$ are the positions of N particles moving in \mathbb{R}^2 , identified with the complex plane, and the constant c_{Lau} is a normalization factor (depending on N and ℓ). For fermions, ℓ is odd and ≥ 3 (the case $\ell = 1$ corresponds to noninteracting fermions), while for bosons $\ell \geq 2$ is even. Bosonic wave functions of this type are potentially relevant for atomic gases in artificial magnetic fields [6, 8, 19, 26, 29, 37]. The analysis below applies to Laughlin states of both symmetry types.

The Laughlin state (1) is a special kind of wave function in the Lowest Landau level (LLL) of a Hamiltonian with a strong magnetic field perpendicular to the plane where the particles move. The general form of such functions is

$$\Psi(z_1, \dots, z_N) = A(z_1, \dots, z_N) e^{-\sum_{i=1}^N |z_i|^2/2} \quad (2)$$

with A analytic, antisymmetric for fermions and symmetric for bosons. Also square integrability is required, but A is not restricted to polynomials. If there is strong repulsive two-body interaction between the particles, the Laughlin function (1) is a natural trial function for low energy states. The factors $(z_i - z_j)^\ell$ indeed suppress interactions by producing strong correlations between the particle positions. For certain zero-range interactions, (1) is even an exact ground state [11, 27, 36].

The success of Laughlin's theory of the FQHE fractions $1/\ell$, based on (1), depends crucially on the fact that the Laughlin wave function behaves as an incompressible liquid, whose response to perturbations and external fields is extremely rigid. This in fact has two distinct aspects:

1. The Laughlin wave function is an approximate ground state for the many-body Hamiltonian, and its energy is separated from the rest of the spectrum by a gap.

2. Modifications of the Laughlin wave function that stay within the ground eigenspace of the many-body Hamiltonian cannot increase the local one-particle density beyond a fixed value.

The main evidence for these facts is experimental and numerical. For model zero-range interactions, the justification of Property **1** amounts to a proof of an N -independent spectral gap, see e.g. [19, Section 2.2] or [32, Section 2.1] and references therein. We are not aware of a solution to this important problem.

Property **2** is not always recognized as a separate issue, but is also crucial for Laughlin's original argument supporting quantization of the Hall conductivity [13, 16] and for FQHE physics in general. Its importance lies in understanding the effect of an external potential that would try to concentrate the density as much as possible in energetically favorable places.

If the energy gap in Property **1** is sufficiently large we can exclude a jump across it and restrict attention to wave functions with the same interaction energy as the Laughlin function. We are thus led to study the set of all normalized wave functions of the form

$$\Psi_F(z_1, \dots, z_N) = F(z_1, \dots, z_N) \Psi_{\text{Lau}}(z_1, \dots, z_N) \quad (3)$$

with F analytic and symmetric under exchange of the z_i . This form exhausts the class of functions that minimize the magnetic kinetic energy and at the same time avoid repulsive interactions by vanishing at least as $(z_i - z_j)^\ell$ as z_i and z_j come together. In the bosonic case and with $\ell = 2$ these are exactly the ground states of the contact interaction [32, Section 2.1]. We shall refer to the class of states of the form (3) as *fully correlated states*.

Physically, (3) can lead to the addition of quasi-holes (zeros of the wave-function) to the Laughlin state, essentially arbitrary correlations between the particles' and quasi-holes' locations being allowed. It is intuitive that this leads to a decrease of the *global* density. It is however far from obvious that no *local* increase of the density can occur anywhere. Several quasi-holes arranged tightly on a circle could, perhaps, increase the density inside the

circle. Moreover, F need not contain any zeros at all like, e.g., $\exp(c \sum_i z_i^2)$, which stretches the support of the density in one direction and shrinks it in another.

In this Letter, we report on recent density bounds for fully correlated states that demonstrate the validity of Property **2** without invoking Property **1**. They hold essentially on length scales $O(N^{1/4})$, much smaller than the full extent of the liquid, which is of order $N^{1/2}$, and are related to earlier partial results proved in [33, 34]. The full proofs are somewhat involved and shall be presented elsewhere [23] but we sketch the main arguments below, and discuss physical applications.

In his pioneering paper [14], Laughlin already argued that the one-particle density of the state (1) has the form of a circular droplet of radius $\sqrt{\ell N}$ where the density takes the constant value $(\pi\ell)^{-1}$. The argument was based on the plasma analogy, where the absolute square of the wave function is written as the Gibbs distribution of a classical 2D Coulomb gas, and subsequently treated by a mean-field approximation.

Recent mathematical analysis has confirmed the validity of this approximation for the Laughlin states and extended it to more general states of the form (3), including Laughlin’s ‘quasi hole’ states where the prefactor F is a product where each factor depends on a single variable [31, 32]. This analysis was generalized to other prefactors of a special kind in [33]. A common feature that emerged was an upper bound on the one-particle density of magnitude $(\pi\ell)^{-1}$, which is the density of the Laughlin state itself. Such a bound was called an *incompressibility estimate* in [33] because it is a manifestation of the resistance of the Laughlin state against attempts to compress its density.

The mean-field methods of [31–33] are not applicable to general prefactors F . The question of a bound on the density for the general case was treated in [34] with an entirely different technique, rooted in 2D potential theory. The bound obtained was, however, four times the expected optimal value $(\pi\ell)^{-1}$. In this Letter we show that an improved version of the potential theoretic method leads to the correct optimal bound for *arbitrary* F .

As indicated by numerical studies [7], the incompressibility bound for the density cannot be expected to hold pointwise for finite N . We prove, however, that it holds for local averages. Denote

$$\rho_F(z) = N \int_{\mathbb{R}^{2(N-1)}} |\Psi_F(z, z_2, \dots, z_N)|^2 dz_2 \dots dz_N \quad (4)$$

the particle density of the state (3). Our main result is

Theorem 1 (Density bound). *For any $\alpha > 1/4$ and any disk D of radius N^α we have*

$$\int_D \rho_F \leq \frac{1}{\pi\ell} |D| (1 + o(1)) \quad (5)$$

where $|D|$ is the area of the disk and $o(1)$ tends to zero as $N \rightarrow \infty$.

Recall that the magnetic length is $1/\sqrt{2}$. The average can thus be taken on “mesoscopic scales”, much smaller than the full extent of the state, $O(N^{1/2})$, but not quite down to the expected finest scale, $O(1)$. A simple argument shows that (5) implies the analogous result for any open set, not just a disk.

It follows that, for a continuous confining potential V , the smallest potential energy obtainable with a fully correlated state

$$E_V(\ell, N) := \min \left\{ \int_{\mathbb{R}^2} V \rho_F, \Psi_F \text{ of the form (3)} \right\} \quad (6)$$

is bounded below by N times the “bathtub energy” [22, Theorem 1.14]

$$E_V^{\text{bt}}(\ell) := \min \left\{ \int_{\mathbb{R}^2} V \rho \mid 0 \leq \rho \leq \frac{1}{\pi\ell}, \int_{\mathbb{R}^2} \rho = N \right\}, \quad (7)$$

i.e. the minimum of the potential energy over all densities ρ bounded by $(\pi\ell)^{-1}$. This lower bound

$$E_V(\ell, N) \gtrsim N E_V^{\text{bt}}(\ell) \quad (8)$$

holds for large N provided V varies only on scales much larger than $N^{1/4}$.

The bound (5) means that any compression of the particle density above the “magic value” $(\pi\ell)^{-1}$ that one could imagine to accommodate the variations of an external potential would make us leave the class of fully correlated states, with corresponding increase in either the magnetic kinetic energy or the interaction energy. Assuming the spectral gap mentioned in Property **1**, no such density bump is allowed. This justifies two things *a posteriori*:

- That it is legitimate to neglect disorder in the sample and/or small external electric fields, as is done as a first approximation in the derivation of FQHE wave-functions.
- Laughlin’s argument [13, 15] (see also [12, Sections 4.4, 9.3 and 9.5]) that switching on an electric current moves electrons transversally without creating any charge accumulation, and generates a Hall conductivity of value $1/\ell$.

It has been proposed (see [6, 8, 37] for reviews) that Laughlin wave functions could be created in cold atomic gases, either by rapid rotation or by applying artificial magnetic fields. In this context, a magneto-optical trap confines the gas. Some recent proposals to reach the Laughlin state [26, 29] involve some non-trivial engineering of the trapping potential. How the Laughlin state responds is, therefore, of importance for the experimental set-up.

In addition, the precursor of FQHE states in a rapidly rotating Bose gas is a Bose-Einstein condensate (see [19, 25] and references therein). Observing the distinctively flat profile of the Laughlin state would already be a strong indication of the transition to the FQHE regime. A more

complete probe could be the response of the gas to variations of the trapping potential: the Bose condensate follows the trap by taking a Thomas-Fermi-like shape (see [1, 2, 5] and references therein). The Laughlin state essentially does not respond to such variations, as exemplified by our main theorem.

We also point out that a combination of Theorem 1 with estimates obtained in [31, 32] leads to the following improvement of [33, Corollary 2.3]:

Corollary 2 (Optimization of the energy in radial traps). *Let $V(x) = |x|^s$ with $s > 0$. Then the potential energy within the class (3) is minimized by the Laughlin state ($F = 1$):*

$$\lim_{N \rightarrow \infty} \frac{\int_{\mathbb{R}^2} V \rho_1}{E_V(\ell, N)} = \lim_{N \rightarrow \infty} \frac{N E_V^{\text{bt}}(\ell)}{E_V(\ell, N)} = 1 \quad (9)$$

where $E_V(N, \ell)$ and $E_V^{\text{bt}}(\ell)$ are defined in (6) and (7) respectively.

It is remarkable that the Laughlin state stays an approximate minimizer in any power-law trap (the result actually holds for more general radial increasing potentials). No matter how steep and narrow a potential well one imposes, it is impossible to compress the Laughlin state while keeping the form (3), i.e., without jumping across the spectral gap. Extensions of Corollary 2 to mexican-hat-like traps as in [33, Corollary 2.3] are also possible.

We now turn to sketching the proof of Theorem 1, postponing the details to a planned longer paper. It is convenient to change variables and consider the scaled N -particle probability density

$$\mu_F(Z_N) := N^N \left| \Psi_F \left(\sqrt{N} Z_N \right) \right|^2 \quad (10)$$

corresponding to the wave-function (3). This has an extension $O(1)$ for the Laughlin state, $F = 1$. Here Z_N stands for (z_1, \dots, z_N) . The scaled 1-particle probability density is

$$\begin{aligned} \mu_F^{(1)}(z) &= \int_{\mathbb{R}^{2(N-1)}} \mu_F(z, z_2, \dots, z_N) dz_2 \dots dz_N \\ &= \rho_F \left(\sqrt{N} z \right) \end{aligned} \quad (11)$$

The first step is to write the N -particle density as a Gibbs factor (Laughlin's plasma analogy),

$$\mu_F(Z_N) = \mathcal{Z}_N^{-1} \exp \left(-\frac{1}{T} H_N(Z_N) \right), \quad (12)$$

with effective "temperature" $T = N^{-1}$ and the Hamiltonian

$$H_N(Z_N) = \sum_{j=1}^N |z_j|^2 - \frac{2\ell}{N} \sum_{i < j} \log |z_i - z_j| + W_N(Z_N) \quad (13)$$

where

$$W_N(Z_N) = -\frac{2}{N} \log \left| F \left(\sqrt{N} Z_N \right) \right|. \quad (14)$$

The term $W_N(Z_N)$ has the important property of being superharmonic in each variable:

$$-\nabla_{z_i}^2 W_N(Z_N) \geq 0 \text{ for all } i. \quad (15)$$

This holds because F is analytic and is, in fact, the only property of W_N that is used in our method.

A precursor of the desired bound (5) for ρ_F is the fact that the local density of points in a minimizing configuration for $H_N(Z_N)$ is everywhere bounded above by $N(\pi\ell)^{-1}$ for large N . This is the core of the proof of the theorem, and a signature of screening properties of the effective plasma. To establish it, we introduce and study an auxiliary minimization problem, which is, mathematically, a cousin of the Thomas-fermi energy minimization problem for molecules [21].

For K fixed points $x_i \in \mathbb{R}^2$ ("nuclei") we define an energy for functions σ on \mathbb{R}^2 ("electron density") by

$$\mathcal{E}^{\text{TF}}[\sigma] = - \int_{\mathbb{R}^2} V_{\text{nuc}}(x) \sigma(x) dx + D(\sigma, \sigma) \quad (16)$$

with

$$V_{\text{nuc}}(x) = - \sum_{i=1}^K \log |x - x_i| \quad (17)$$

and

$$D(\sigma, \sigma') = -\frac{1}{2} \iint_{\mathbb{R}^2 \times \mathbb{R}^2} \sigma(x) \log |x - x'| \sigma'(x') dx dx'. \quad (18)$$

This functional is minimized under the subsidiary conditions

$$0 \leq \sigma \leq 1, \quad \int_{\mathbb{R}^2} \sigma = K. \quad (19)$$

In physical terms, this model describes a neutral 2D molecule consisting of fixed nuclei and mobile electrons, with Coulomb interactions. The interpretation of the constraint $0 \leq \sigma \leq 1$ is that the kinetic energy of the electrons is zero for densities ≤ 1 and ∞ for densities > 1 .

The basic facts about this TF model are: (1) There exists a unique minimizer, σ^{TF} . (2) The minimizer has compact support. (3) Apart of a set of measure zero, σ^{TF} takes only the values 0 or 1. (4) The Thomas-Fermi equation holds:

$$\Phi^{\text{TF}}(x) = \begin{cases} \geq 0 & \text{if } \sigma^{\text{TF}}(x) = 1 \\ 0 & \text{if } \sigma^{\text{TF}}(x) = 0 \end{cases} \quad (20)$$

where

$$\Phi^{\text{TF}}(x) = V_{\text{nuc}}(x) + \int_{\mathbb{R}^2} \log |x - x'| \sigma^{\text{TF}}(x') dx'$$

is the total electrostatic potential of the molecule.

According to the TF equation the support of σ^{TF} is the same as the support of the potential Φ^{TF} , which is continuous away from the “nuclei”. Denote by $\Sigma^{\text{TF}}(x_1, \dots, x_K)$ the open set where Φ^{TF} is strictly larger than 0. Some important properties are:

(1) The area of $\Sigma^{\text{TF}}(x_1, \dots, x_K)$ is equal to K . (2) $\Sigma^{\text{TF}}(x_1, \dots, x_{K-1}) \subset \Sigma^{\text{TF}}(x_1, \dots, x_K)$. (3) For a single nucleus at x_1 , $\Sigma^{\text{TF}}(x_1)$ is the disc with center x_1 and radius $\pi^{-1/2}$.

Consider now a scaled version of (13),

$$\mathcal{H}(X_N) = \frac{\pi}{2} \sum_{i=1}^N |x_i|^2 - \sum_{1 \leq i < j \leq N} \log |x_i - x_j| + \mathcal{W}(X_N), \quad (21)$$

with \mathcal{W} symmetric and superharmonic in each variable $x_i \in \mathbb{R}^2$ and $X_N = (x_1, \dots, x_N)$. A key property of minimizing configurations of \mathcal{H} is stated in the following:

Lemma 3 (Exclusion rule). *Let $X_N^0 = \{x_1^0, \dots, x_N^0\}$ be a minimizing configuration of points for \mathcal{H} . For any subset of points $y_1, \dots, y_K, y_{K+1} \in X_N^0$,*

$$y_{K+1} \notin \Sigma^{\text{TF}}(y_1, \dots, y_K). \quad (22)$$

The proof of this intermediary result is sketched in the Appendix. To make the result plausible, observe the following: the first term in (21) is the electrostatic potential generated by a constant background of charge density -1 . Using Equation (20), the electrostatic potential generated by the points y_1, \dots, y_K is completely screened by the part of the background potential generated by the region $\Sigma^{\text{TF}}(y_1, \dots, y_K)$. What the result says is that no other point of a minimizing configuration can lie inside this screening region. The reason is two-fold:

- if another point y_{K+1} lay inside the screening region, one could decrease the sum of the first two terms in (21) by moving y_{K+1} to any position on the boundary.
- the last term \mathcal{W} in (21), being superharmonic, is generated by a positive charge distribution. One can thus always decrease \mathcal{W} by moving y_{K+1} to some point on the boundary. Such a move decreases at the same time the sum of the first two terms.

The particular case $K = 1$ of Lemma 3 goes back to an unpublished theorem of Lieb, used in [34]: The minimal distance between points in a minimizing configuration of \mathcal{H} is not less than $1/\sqrt{\pi}$. This property shows that the density of points is in any case bounded above by 4. The general exclusion rule for all K implies more. The density is, in fact, asymptotically bounded above by 1:

Lemma 4 (Exclusion rule implies density bound). *For $R > 0$ let $n(R)$ denote the maximum number of any points $\{y_1, \dots, y_n\}$ that a disk $D(R)$ of radius R can accommodate while respecting the exclusion rule (22). Then*

$$\limsup_{R \rightarrow \infty} \frac{n(R)}{\pi R^2} \leq 1. \quad (23)$$

We present a proof sketch in the Appendix. The main point is that, if a large region contains a density of points > 1 we can reach a contradiction by considering two facts:

- the potential Φ^{TF} generated by the points y_i contained in $D(R)$ and the corresponding exclusion set $\Sigma^{\text{TF}}(y_1, \dots, y_n)$ must vanish at all the points y_j lying outside of $D(R)$, by the exclusion rule and (20). This leads to a uniform upper bound on Φ^{TF} outside of $D(R)$.
- the same potential is generated by an overall positive charge density, because the nuclear charge in $D(R)$, which is $> \pi R^2$ by assumption, is not fully screened by the part of the negative charge density σ^{TF} lying in $D(R)$, at most equal to the area πR^2 because $\sigma^{\text{TF}} \leq 1$. A lower bound on the circular average of Φ^{TF} outside of $D(R)$ follows.

It turns out that, for R large enough, the two estimates obtained in this way contradict one another.

After scaling, $x \rightarrow z = \sqrt{\frac{\pi \ell}{N}} x$, Lemma 2 applies to the Hamiltonian (13) and implies that in any minimizing configuration $\{z_1^0, \dots, z_N^0\}$ of (13) the number of points z_i^0 contained in any disc of radius $R \gg N^{-1/2}$ is no larger than $N(\pi \ell)^{-1}$ times the area of the disc. This is the gist of the proof. From there, two main steps are left to conclude the proof of Theorem 1, see the Appendix.

We point out that our density upper bound holds down to the finest possible scale for ground states of the plasma Hamiltonian (13), i.e. on length scales $\gg N^{-1/2}$ (see [30] where the corresponding lower bound is proved in the purely Coulombic case $W = 0$). When applying this result to Gibbs states of (13) we have to restrict to length scales $\gg N^{-1/4}$, but this is likely to be due to a technical limitation of our method. It was indeed recently proved that, for the purely Coulombic Hamiltonian, the expected microscopic density estimate holds for low temperature Gibbs states [3, 4, 17, 18] (see also [28] for ground states of higher dimensional Coulomb and Riesz gases). It remains to be seen whether a combination of our methods with those of [3, 17] could improve our results. Another question for future investigations is the extension of Corollary 2 to more general potentials V .

Conclusion: We have considered perturbations of the Laughlin state that may arise to accommodate external potentials while keeping the system in the lowest Landau band and preserving the original correlations. We proved rigorously that no such perturbation can raise the particle density anywhere beyond the Laughlin value $1/(\pi \ell)$. Our theorem holds on length scales $\gg N^{1/4}$ which, while large compared to the magnetic length, are microscopic compared to the system’s macroscopic size, $N^{1/2}$.

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Appendix : Elements of proof

Here we sketch, for completeness, the proofs of Lemmas 3 and 4. The arguments are kept short; full details are postponed to a longer paper.

Proof of Lemma 3. By symmetry of the Hamiltonian we may, without loss, choose $y_i = x_i^0, 1 \leq i \leq K + 1$. Consider then fixing all points but x_{K+1}^0 . The energy to consider is then

$$G(x) = \mathcal{H}(x_1^0, \dots, x_K^0, x, x_{K+2}^0, \dots, x_N^0). \quad (24)$$

We claim that if $x \in \Sigma^{\text{TF}}(x_1^0, \dots, x_K^0) \equiv \Sigma^{\text{TF}}$ then there is an $\tilde{x} \in \partial\Sigma^{\text{TF}}$, the boundary of Σ^{TF} , such that $G(\tilde{x}) < G(x)$. Thus the minimizing point x_{K+1}^0 cannot lie in Σ^{TF} .

To prove the claim, we add and subtract a term $-\int_{\Sigma^{\text{TF}}} \log|x - x'|dx'$ to write $G(x) = \Phi(x) + R(x)$ with

$$\Phi(x) = -\sum_{i=1}^K \log|x - x_i^0| + \int_{\Sigma^{\text{TF}}} \log|x - x'|dx' \quad (25)$$

and

$$R(x) = \frac{\pi}{2}|x|^2 - \int_{\Sigma^{\text{TF}}} \log|x - x'|dx' - \sum_{i=K+2}^N \log|x - x_i^0| + W(x) + \text{const}. \quad (26)$$

Now, Φ is precisely the TF potential corresponding to ‘nuclear charges’ at x_i^0, \dots, x_K^0 . Hence, using (20), $\Phi > 0$ on Σ^{TF} and zero on the boundary $\partial\Sigma^{\text{TF}}$. The first two terms in R are harmonic on Σ^{TF} when taken together. (The Laplacian applied to the first term gives 2π and to the second term -2π on Σ^{TF} .) The other terms are superharmonic on Σ^{TF} . Thus, R takes its minimum on the boundary, so there is a $\tilde{x} \in \partial\Sigma^{\text{TF}}$ with $R(x) \geq R(\tilde{x})$. On the other hand, $\Phi(x) > 0 = \Phi(\tilde{x})$ so $G(x) > G(\tilde{x})$. \square

Proof of Lemma 4. The proof is indirect. Assuming that for some $\delta > 0$ there are arbitrary large radii with the property that the disk $D(R)$ contains at least $(1 + \delta)\pi R^2$ points. This leads to a contradiction with (22) and we present here a sketch of the argument. Full details are given in a longer paper.

By taking the maximal δ (which is in any case ≤ 3) we may, without restriction, assume that the density is also at least $(1 + \delta)$ in the annulus \mathcal{A} of width $\delta \cdot R$ around $D(R)$. Since the density is everywhere bounded above by 4, apart from the points $y_i \in D(R), i = 1, \dots, n$ there must be points $y_j \in \mathcal{A}, j = n + 1, \dots, m$ in the configuration such that every point in \mathcal{A} is at most a distance $O(1)$ from one of the y_j . The TF potential Φ^{TF} generated by the y_i and the corresponding exclusion set

$\Sigma^{\text{TF}}(y_1, \dots, y_n)$ must vanish at the y_j by the exclusion rule and (20).

On the other hand, after scaling the variables by R^{-1} and extracting a factor R^2 one can show that the gradient of the TF potential is uniformly bounded in the scaled annulus $R^{-1}\mathcal{A}$. The distance between the scaled y_j is now R^{-1} so the scaled potential goes to zero uniformly on $R^{-1}\mathcal{A}$ as $R \rightarrow \infty$. The same holds then for the circular average of the scaled potential.

We claim, however, that the latter is strictly bounded away from zero close to the radius 1 (corresponding to radius R in the unscaled annulus). This follows from Newton’s theorem, because the nuclear charge in $D(R)$, which is $(1 + \delta)\pi R^2$ by assumption, is not fully screened by the part of the negative charge density σ^{TF} lying in $D(R)$, which is at most equal to the area πR^2 because $\sigma^{\text{TF}} \leq 1$. There is thus a contradiction for R large enough. \square

As indicated in the main text, with Lemmas 3 and 4 at our disposal, the bound (5) of Theorem 1 holds if $\mu_F^{(1)}$ in (11) is replaced by the empirical measure

$$\rho^0(z) = N^{-1} \sum_i \delta(z - z_i^0)$$

of a minimizing configuration $\{z_1^0, \dots, z_n^0\}$ for H_N . The remaining task is to show that this bound holds also for the Gibbs state of H_N . Here we use crucially that the ‘temperature’ T in the latter scales as N^{-1} so that the Gibbs measure is essentially concentrated on ground state configurations for large N . Turning this intuition into a proof follows the lines of [34, Section 3]: To access the 1-particle density we prove free energy upper and lower bounds for a perturbed version of the Hamiltonian. We show that the latter satisfies, approximately, the same bounds as the unperturbed (13) and then translate the free-energy estimates into density estimates using a Feynman-Hellmann-type argument.

Strictly speaking, this strategy requires an a priori bound ensuring that Ψ_F lives on length scales of order $N^{1/2}$, e.g.

$$\langle \Psi_F, L_N \Psi_F \rangle \leq CN^2 \quad (27)$$

where L_N is the total angular momentum operator and C is independent of N and F . This assumption is very reasonable physically, and can in fact be eliminated using a suitable localization procedure.

Finally, the extension of Theorem 1 from disks to arbitrary open sets uses the ‘cheese theorem’ (see [24, Section 14.4] or [20, Theorem 14]) which asserts that any reasonable open set can be efficiently covered with disks, up to a residual set of arbitrarily small area.

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- [1] A. AFTALION, X. BLANC, AND J. DALIBARD, *Vortex patterns in a fast rotating Bose-Einstein condensate*, Phys. Rev. A, 71 (2005), p. 023611.
- [2] A. AFTALION, X. BLANC, AND F. NIER, *Vortex distribution in the lowest Landau level*, Phys. Rev. A, 73 (2006), p. 011601(R).
- [3] R. BAUERSCHMIDT, P. BOURGADE, M. NIKULA, AND H.-T. YAU, *Local density for two-dimensional one-component plasma*. arXiv:1510.02074, 2015.
- [4] ———, *The two-dimensional Coulomb plasma: quasi-free approximation and central limit theorem*. arXiv:1609.08582, 2016.
- [5] X. BLANC AND N. ROUGERIE, *Lowest-Landau-Level vortex structure of a Bose-Einstein condensate rotating in a harmonic plus quartic trap*, Phys. Rev. A, 77 (2008), p. 053615.
- [6] I. BLOCH, J. DALIBARD, AND W. ZWERGER, *Many-body physics with ultracold gases*, Rev. Mod. Phys., 80 (2008), pp. 885–964.
- [7] O. CIFTJÁ, *Monte Carlo study of Bose Laughlin wave function for filling factors 1/2, 1/4 and 1/6*, Europhys. Lett., 74 (2006), pp. 486–492.
- [8] N. R. COOPER, *Rapidly rotating atomic gases*, Advances in Physics, 57 (2008), pp. 539–616.
- [9] S. GIRVIN, *Introduction to the fractional quantum Hall effect*, Séminaire Poincaré, 2 (2004), pp. 54–74.
- [10] M. O. GOERBIG, *Quantum Hall effects*. arXiv:0909.1998, 2009.
- [11] F. D. M. HALDANE, *Fractional quantization of the Hall effect: A hierarchy of incompressible quantum fluid states*, Phys. Rev. Lett., 51 (1983), pp. 605–608.
- [12] J. K. JAIN, *Composite fermions*, Cambridge University Press, 2007.
- [13] R. B. LAUGHLIN, *Quantized Hall conductivity in two dimensions*, Phys. Rev. B, 23 (1981), p. 5632.
- [14] ———, *Anomalous quantum Hall effect: An incompressible quantum fluid with fractionally charged excitations*, Phys. Rev. Lett., 50 (1983), pp. 1395–1398.
- [15] ———, *Elementary theory : the incompressible quantum fluid*, in The quantum Hall effect, R. E. Prange and S. E. Girvin, eds., Springer, Heidelberg, 1987.
- [16] ———, *Nobel lecture: Fractional quantization*, Rev. Mod. Phys., 71 (1999), pp. 863–874.
- [17] T. LEBLÉ, *Local microscopic behavior for 2D Coulomb gases*. arXiv:1510.01506, 2015.
- [18] T. LEBLÉ AND S. SERFATY, *Fluctuations of two-dimensional Coulomb gases*. arXiv:1609.08088, 2016.
- [19] M. LEWIN AND R. SEIRINGER, *Strongly correlated phases in rapidly rotating Bose gases*, J. Stat. Phys., 137 (2009), pp. 1040–1062.
- [20] E. H. LIEB, *The stability of matter*, Rev. Mod. Phys., 48 (1976), pp. 553–569.
- [21] E. H. LIEB, *Thomas-Fermi and related theories of atoms and molecules*, Rev. Mod. Phys., 53 (1981), pp. 603–641.
- [22] E. H. LIEB AND M. LOSS, *Analysis*, vol. 14 of Graduate Studies in Mathematics, American Mathematical Society, Providence, RI, 2nd ed., 2001.
- [23] E. H. LIEB, N. ROUGERIE, AND J. YNGVASON. In preparation.
- [24] E. H. LIEB AND R. SEIRINGER, *The Stability of Matter in Quantum Mechanics*, Cambridge Univ. Press, 2010.
- [25] E. H. LIEB, R. SEIRINGER, AND J. YNGVASON, *Yrast line of a rapidly rotating Bose gas: Gross-Pitaevskii regime*, Phys. Rev. A, 79 (2009), p. 063626.
- [26] A. MORRIS AND D. FEDER, *Gaussian potentials facilitate access to quantum Hall states in rotating Bose gases*, Phys. Rev. Lett., 99 (2007), p. 240401.
- [27] T. PAPANBROCK AND G. F. BERTSCH, *Rotational spectra of weakly interacting Bose-Einstein condensates*, Phys. Rev. A, 63 (2001), p. 023616.
- [28] M. PETRACHE AND S. ROTA-NODARI, *Equidistribution of jellium energy for Coulomb and Riesz interactions*. arXiv:1609.03849, 2016.
- [29] M. RONCAGLIA, M. RIZZI, AND J. DALIBARD, *From rotating atomic rings to quantum Hall states*, www.nature.com, Scientific Reports, 1 (2011).
- [30] S. ROTA NODARI AND S. SERFATY, *Renormalized energy equidistribution and local charge balance in 2d Coulomb systems*, Int. Math. Res. Not., 11 (2015), pp. 3035–3093.
- [31] N. ROUGERIE, S. SERFATY, AND J. YNGVASON, *Quantum Hall states of bosons in rotating anharmonic traps*, Phys. Rev. A, 87 (2013), p. 023618.
- [32] ———, *Quantum Hall phases and plasma analogy in rotating trapped Bose gases*, J. Stat. Phys., 154 (2014), pp. 2–50.
- [33] N. ROUGERIE AND J. YNGVASON, *Incompressibility estimates for the Laughlin phase*, Comm. Math. Phys., 336 (2015), pp. 1109–1140.
- [34] ———, *Incompressibility estimates for the Laughlin phase, part II*, Comm. Math. Phys., 339 (2015), pp. 263–277.
- [35] H. STÖRMER, D. TSUI, AND A. GOSSARD, *The fractional quantum Hall effect*, Rev. Mod. Phys., 71 (1999), pp. S298–S305.
- [36] S. TRUGMAN AND S. KIVELSON, *Exact results for the fractional quantum Hall effect with general interactions*, Phys. Rev. B, 31 (1985), p. 5280.
- [37] S. VIEFERS, *Quantum Hall physics in rotating Bose-Einstein condensates*, J. Phys. C, 20 (2008), p. 123202.