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Phys. Rev. Lett. **107**, 040501 — Published 20 July 2011

DOI: [10.1103/PhysRevLett.107.040501](https://doi.org/10.1103/PhysRevLett.107.040501)

# Computational Difficulty of Computing the Density of States

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We study the computational difficulty of computing the ground state degeneracy and the density of states for local Hamiltonians. We show that the difficulty of both problems is exactly captured by a class which we call  $\#BQP$ , which is the counting version of the quantum complexity class  $QMA$ . We show that  $\#BQP$  is not harder than its classical counting counterpart  $\#P$ , which in turn implies that computing the ground state degeneracy or the density of states for classical Hamiltonians is just as hard as it is for quantum Hamiltonians.

Understanding the physical properties of correlated quantum many-body systems is a problem of central importance in condensed matter physics. The density of states, defined as the number of energy eigenstates per energy interval, plays a particularly crucial role in this endeavor. It is a key ingredient when deriving many thermodynamic properties from microscopic models, including specific heat capacity, thermal conductivity, band structure, and (near the Fermi energy) most electronic properties of metals. Computing the density of states can be a daunting task however, as it in principle involves diagonalizing a Hamiltonian acting on an exponentially large space, though other more efficient approaches which might take advantage of the structure of a given problem are not a priori ruled out.

In this Letter, we precisely quantify the difficulty of computing the density of states by using the powerful tools of quantum complexity theory. Quantum complexity aims at generalizing the well-established field of classical complexity theory to assess the difficulty of tasks related to quantum mechanical problems, concerning both the classical difficulty of simulating quantum systems as well as the fundamental limits to the power of quantum computers. In particular, quantum complexity theory has managed to explain the difficulty of computing ground state properties of quantum spin systems in various settings, such as two-dimensional (2D) lattices [1] and even one-dimensional (1D) chains [2], as well as fermionic systems [3].

We will determine the computational difficulty of two problems: First, computing the density of states of a local Hamiltonian, and second, counting the ground state degeneracy of a local gapped Hamiltonian; in both cases, the result holds even if the Hamiltonian is restricted to act on a 2D lattice of qubits, or on a 1D chain. To this end, we will introduce the quantum counting class  $\#BQP$  (sharp BQP), which constitutes the natural counting version of the class  $QMA$  (Quantum Merlin Arthur) which itself captures the difficulty of computing the ground state energy of a local Hamiltonian [4, 5]. Vaguely speaking,  $\#BQP$  counts the number of possible “quantum solutions” to a quantum problem that can be verified using a

quantum computer. We show that both problems, computing the density of states and counting the ground state degeneracy, are complete problems for the class  $\#BQP$ , i.e., they are among the hardest problems in this class.

Having quantified the difficulty of computing the density of states and counting the number of ground states, we proceed to relate  $\#BQP$  to known classical counting complexity classes, and show that  $\#BQP$  equals  $\#P$  (under weakly parsimonious reductions). Here, the complexity class  $\#P$  counts the number of satisfying assignments to any efficiently computable boolean function. This can be a very hard problem which is believed to take exponential time; in particular, it is at least as hard as deciding whether the function has at least one satisfying input, i.e., the complexity class  $NP$ . Examples for  $\#P$ -complete problems (i.e., the hardest problems in that class) include counting the number of colorings of a graph, or computing the permanent of a matrix with binary entries. Phrased in terms of Hamiltonians, what we show is that computing the density of states and counting the ground state degeneracy of a classical spin system is just as hard as solving the same problem for a quantum Hamiltonian.

*Quantum complexity classes.*—Let us start by introducing the relevant complexity classes. The central role in the following is taken by the *verifier*  $V$ , which verifies “quantum solutions” (also called *proofs*) to a given problem. More formally, a verifier checking an  $n$ -qubit quantum proof (that is, a quantum state  $|\psi\rangle$ ) consists of a  $T = \text{poly}(n)$  length quantum circuit  $U = U_T \cdots U_1$  (with local gates  $U_i$ ) acting on  $m = \text{poly}(n)$  qubits, which takes the  $n$ -qubit quantum state  $|\psi\rangle_I$  as an input, together with  $m - n$  initialized ancillas,  $|0\rangle_A \equiv |0 \cdots 0\rangle_A$ , applies  $U$ , and finally measures the first qubit in the  $\{|0\rangle_1, |1\rangle_1\}$  basis to return 1 (“proof accepted”) or 0 (“proof rejected”). Then, the class  $QMA$  contains all problems of the form: “Decide whether there exists a  $|\psi\rangle$  such that  $p_{\text{acc}}(V(\psi)) > a$ , or whether  $p_{\text{acc}}(V(\psi)) < b$  for all  $|\psi\rangle$ , for some chosen  $a - b > 1/\text{poly}(n)$ , given that one is the case”. Here, the acceptance probability of a state  $|\psi\rangle$  is  $p_{\text{acc}}(V(\psi)) := \langle \psi | \Omega | \psi \rangle$ , with

$$\Omega = (\mathbb{1}_I \otimes \langle 0 |_A) U^\dagger (|1\rangle_1 \langle 1|_1 \otimes \mathbb{1}) U (\mathbb{1}_I \otimes |0\rangle_A), \quad (1)$$

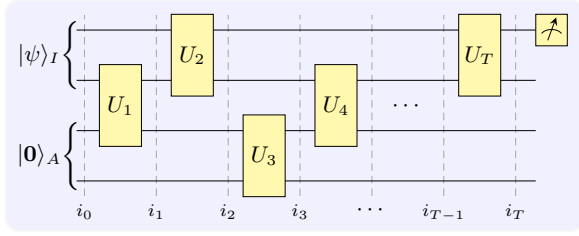


FIG. 1: A QMA verifier consists of a sequence of  $T$  local unitary gates acting on the “quantum proof”  $|\psi\rangle$  and an ancillary register initialized to  $|0\rangle$ . The final measurement on the first qubit returns  $|1\rangle$  or  $|0\rangle$  to accept or reject the proof, respectively. Transition probabilities can be computed by doing a “path integral” over all intermediate configurations  $(i_k)_k$ .

which we illustrate in Fig. 1.

The idea behind this definition is that QMA quantifies the difficulty of computing the ground state energy  $E_0(H)$  of a local Hamiltonian  $H$  up to  $1/\text{poly}(n)$  accuracy. Let the verifier be a circuit estimating  $\langle\psi|H|\psi\rangle$ ; then a black box solving QMA problems can be used to compute  $E_0(H)$  up to  $1/\text{poly}(n)$  accuracy by binary search using a single QMA query. Note also that QMA is the quantum version of the class NP, where one is given an efficiently computable boolean function  $f(x) \in \{0, 1\}$  and one must figure out if there is an  $x$  such that  $f(x) = 1$ .

The class NP has a natural counting version, known as #P. Here, the task is to determine the *number* rather than the existence of satisfying inputs, i.e., to compute  $|\{x : f(x) = 1\}|$ . We will now analogously define #BQP, the counting version of QMA. Consider the verifying map  $\Omega$  of Eq. (1) for a QMA problem, with the additional promise that  $\Omega$  does not have eigenvalues between  $a$  and  $b$ ,  $a - b > 1/\text{poly}(n)$ . Then the class #BQP consists of all problems of the form “compute the dimension of the space spanned by all eigenvectors with eigenvalues  $\geq a$ ”.

An equivalent definition for #BQP (cf. also [6, 7]) is the following: Consider a verifier  $\Omega$  with the additional promise that there exist subspaces  $\mathcal{A} \oplus \mathcal{R} = \mathbb{C}^{2^n}$  such that  $\langle\psi|\Omega|\psi\rangle \geq a$  for all  $|\psi\rangle \in \mathcal{A}$ , and  $\langle\psi|\Omega|\psi\rangle \leq b$  for all  $|\psi\rangle \in \mathcal{R}$ , where again  $a - b > 1/\text{poly}(n)$  – we can think of  $\mathcal{A}$  and  $\mathcal{R}$  as containing the good and bad witnesses, respectively. (Note that there will always be “mediocre” witnesses – the question is whether there exists a decomposition into a good and a bad witness space.) Then, #BQP consists of all problems of the form “compute  $\dim \mathcal{A}$ ”. This number is well-defined, i.e., independent of the choice of  $\mathcal{A}$  and  $\mathcal{R}$ , and moreover, one can easily show that it is equivalent to the definition above, cf. the Supplementary Material.

The gap promise we impose on the spectrum of  $\Omega$  is not present in the definition of QMA (though similarly restricted versions of QMA were defined in [6, 7]). Nevertheless, this promise emerges naturally when considering the counting version: QMA captures the difficulty of

determining the existence of an input state with acceptance probability above  $a$ , up to a “grace interval”  $[b, a]$  in which mistakes are tolerated (i.e., if the largest eigenvalue of  $\Omega$  is in  $[b, a]$ , the oracle can return either outcome). Correspondingly, #BQP captures the difficulty of counting the number of eigenvalues above  $a$ , where eigenvalues in the grace interval  $[b, a]$  can be miscounted. The reason why we choose to define #BQP with a gap promise rather than with a grace interval is the same as for QMA, namely to have a unique outcome associated with any input.

Similarly, the idea of the Hamiltonian formulation of the problem which we will discuss below is to ask for the number of eigenstates in a certain energy interval, where states which are in some small  $1/\text{poly}(n)$  neighborhood of this interval may be miscounted; again, for reasons of rigor we choose to consider only Hamiltonians with no eigenstates in that interval. It should be noted, however, that all of the equivalence proofs we give equally apply if we choose to allow for miscounting of states in those grace intervals instead of requiring them to be empty, as the proofs do not make use of the gap promise itself, but rather show that all states outside those grace intervals are mapped (and thus counted) correctly. Thus, while the actual number returned by the grace interval formulation of the counting problems might change under those mappings due to different treatment of states in the grace interval, it will still be in the correct range.

The class #BQP inherits the important property from QMA of being stable under amplification, that is, the definition of #BQP is not sensitive to the choice of  $a$  and  $b$ . In particular, any  $a - b > 1/\text{poly}(n)$  can be amplified (by building a new poly-size  $\Omega'$  from  $\Omega$ ) such that  $a' = 1 - \exp(-\text{poly}(n))$ ,  $b' = \exp(-\text{poly}(n))$ , and keeping the eigenvalue gap between  $a'$  and  $b'$ , by using a construction called *strong amplification*, cf. Ref. [8]; as shown there, strong amplification acts on all eigenvalues independently and thus also applies to #BQP. The crucial point is that strong amplification works without changing the proof itself, compared to weak amplification which takes multiple copies of the proof as an input. While this is fine for QMA, it does change the dimension of the accepting subspace in an unpredictable way and is thus not an option for the amplification of #BQP.

*Complexity of computing the density of states.*—Let us now show why the class #BQP is relevant for physical applications. In particular, we are going to show that computing the density of states of a local  $n$ -spin Hamiltonian  $H = \sum_i H_i$  with few-body terms  $H_i$ ,  $\|H_i\| \leq 1$ , up to accuracy  $1/\text{poly}(n)$ , is a problem which is complete for #BQP, i.e., it is as hard as any problem in #BQP can be. The same holds true for the (a priori weaker) problem of counting the ground state degeneracy of a local Hamiltonian, given a  $1/\text{poly}(n)$  spectral gap above (note that Bravyi *et al.* [9] suggested this as a definition for a quantum counting class). We can impose additional

restrictions on the interaction structure of our Hamiltonian, and as we will see, the hardness is preserved even for 2D lattices of qubits, or 1D systems.

The problem DOS (density of states) is defined as follows: Given a local Hamiltonian  $H = \sum_i H_i$ , compute the number of orthogonal eigenstates with eigenvalues in an interval  $[E_1, E_2]$  with  $E_2 - E_1 > 1/\text{poly}(n)$ , where we do not allow for eigenvalues within a small grace interval of width  $\Delta = (E_2 - E_1)/\text{poly}(n)$  centered around  $E_1$  and  $E_2$ ; alternatively, we can allow for erroneous counts of eigenstates in that interval. Second, the problem #LH (sharp local Hamiltonian) corresponds to counting the number of ground states of a local Hamiltonian which has a spectral gap  $\Delta = 1/\text{poly}(n)$  above the ground state subspace, given we are told the ground state energy, and where we allow for a small splitting in the ground state energies; again, we can alternatively allow to miscount states in the grace interval.

Clearly, #LH is a special instance of DOS, i.e., solving #LH can be reduced to solving DOS. In order to show that DOS is contained in #BQP, we can use a phase estimation circuit [10] to estimate the energy of any given input  $|\psi\rangle$  and only accept if its energy  $\langle\psi|H|\psi\rangle$  is in the interval  $[E_1, E_2]$ ; as the desired accuracy  $\Delta = 1/\text{poly}(n)$ , this can be done efficiently. A detailed proof (using a more elementary circuit) is given in the Supplementary Material.

Let us now conversely show that #LH is a hard problem for #BQP, that is, any problem in #BQP can be reduced to counting the ground states of some gapped local Hamiltonian [18]. As in turn #LH can be reduced to DOS, which is contained in #BQP, this proves that both #LH and DOS are complete problems for #BQP, i.e., they capture the full difficulty of this class. To this end, we need to start from an arbitrary verifier circuit  $U = U_T \cdots U_1$  and construct a Hamiltonian which has as many ground states as the circuit has accepting inputs (corresponding to the outcome  $|1\rangle_1$  on the first qubit). Let  $\mathcal{A}$  and  $\mathcal{R}$  be the eigenspaces of  $\Omega$  [Eq. (1)] with eigenvalues  $\geq a = 1 - 2^{-\text{poly}(n)}$  and  $\leq b = 2^{-\text{poly}(n)}$ , respectively, and define  $U[\mathcal{R}] := \{U|\psi\rangle_I|0\rangle_A : |\psi\rangle_I \in \mathcal{R}\}$ .

We will follow Kitaev's original construction for a Hamiltonian encoding a QMA verifier circuit [4, 5], which for any proof  $|\psi\rangle_I \in \mathcal{A}$  has the “proof history”  $|\Phi\rangle = \sum_{t=0}^T U_t \cdots U_1 |\psi\rangle_I |0\rangle_A |t\rangle_T$  as its ground state, where the third register is used as a clock. The Hamiltonian  $H = H_{\text{init}} + \sum_{t=1}^T H_{\text{evol}}(t) + H_{\text{final}}$  has three types of terms:  $H_{\text{init}} = \mathbb{1} \otimes (\mathbb{1} - |0\rangle\langle 0|_A) \otimes |0\rangle\langle 0|_T$  makes sure the ancilla is initialized,  $H_{\text{evol}}(t) = -U_t \otimes |t\rangle\langle t-1|_T + \text{h.c.}$  ensures proper evolution from  $t-1$  to  $t$ , and  $H_{\text{final}} = \Pi_{U[\mathcal{R}]} \otimes |T\rangle\langle T|_T$  gives an energy penalty to states  $|\Phi\rangle$  built from proofs  $|\psi\rangle_I \in \mathcal{R}$ . Note that our  $H_{\text{final}}$  differs from the usual choice  $|0\rangle\langle 0|_1 \otimes \mathbb{1} \otimes |T\rangle\langle T|_T$  and is in fact non-local; as we show in the Supplementary Material, this does not significantly change the relevant spectral properties (in particular, we keep the  $1/\text{poly}(n)$  gap, and

the ground state subspace is split at most exponentially). With this choice of  $H_{\text{final}}$ ,  $H$  acts independently on the subspaces spanned by  $\{U_t \cdots U_1 |\psi\rangle_I |x\rangle_A |t\rangle_T\}_{t=0, \dots, T}$  for any  $|\psi\rangle \in \mathcal{A}$  or  $|\psi\rangle \in \mathcal{R}$ , and  $|x\rangle_A$  the computational basis, and the restriction of  $H$  to any of these subspaces describes a random walk which is characterized by the choice of  $|\psi\rangle_I$  and the number of 1's in  $|x\rangle_A$ . These cases can be analyzed independently (see Supplementary Material), and it follows that  $H$  has a  $\dim \mathcal{A}$ -fold degenerate ground state space with a  $1/\text{poly}(n)$  gap above, proving #BQP-hardness of #LH.

This shows that #LH is #BQP-hard for a Hamiltonian which is a sum of  $\log T$ -local terms (i.e., each term acts on  $\log T$  sites), as the clock register is of size  $T$ . In order to obtain a  $k$ -body Hamiltonian, Kitaev suggested to use a unary encoding of the clock (i.e.,  $|t\rangle_T$  is encoded as  $|1 \cdots 10 \cdots 0\rangle$ , with  $t$  1's), so that each Hamiltonian term only acts on three qubits of the clock. However, this makes the Hilbert space of the clock too big, and terms need to be added to the Hamiltonian to penalize illegal clock configurations. These terms divide the Hilbert space into two parts,  $\mathcal{H}_{\text{legal}} \oplus \mathcal{H}_{\text{illegal}}$ . Here,  $\mathcal{H}_{\text{legal}}$  contains only legal clock states, whereas  $\mathcal{H}_{\text{illegal}}$  contains only configurations with illegal clock states [4, 5]. Since no Hamiltonian term couples these two subspaces, the Hamiltonian can be analyzed independently on the two subspaces. It turns out that its restriction to  $\mathcal{H}_{\text{illegal}}$  has an at least  $1/\text{poly}(n)$  higher energy, while on  $\mathcal{H}_{\text{legal}}$ , the Hamiltonian is exactly the same as before. Thus, one finds that the new Hamiltonian still has the right number of ground states, and a  $1/\text{poly}(n)$  spectral gap. The very same argument applies in the case of 1D Hamiltonians, using the QMA-construction of Ref. [2]: Again, the Hamiltonian acts independently on a “legal” and an “illegal” subspace, where the latter has a polynomially larger energy, and the former reproduces the (low-energy) spectrum of the original Hamiltonian [11].

An alternative way to prove QMA-hardness on restricted interaction graphs is to use so-called *perturbation gadgets*, which yield the Hamiltonian of the Kitaev construction above from a perturbative expansion; in particular, this way one can show QMA-hardness of Pauli-type nearest-neighbor Hamiltonians on a 2D square lattice of qubits [1]. As shown in Ref. [12], such gadgets do in fact approximately preserve the whole low-energy part of the spectrum, and thus, our #BQP-hardness proof for #LH still applies to these classes of Hamiltonians. It should be noted, however, that since the eigenvalues are only preserved up to an error  $1/\text{poly}(n)$ , the splitting of the ground state space will now be of order  $1/\text{poly}(n)$ ; however, it can still be chosen to be polynomially smaller than the spectral gap.

*Quantum vs. classical counting complexity.*—As we have seen, the quantum counting class #BQP exactly captures the difficulty of counting the degeneracy of ground states and computing the density of states of local

quantum Hamiltonians. In the following, we will relate  $\#BQP$  to classical counting classes and prove that  $\#BQP$  is equal to  $\#P$ , counting the number of satisfying inputs to a boolean function [19]. In physical terms, this shows that counting the number of ground states or determining the density of states for a quantum Hamiltonian is not harder than either problem is for a classical Hamiltonian.

Clearly,  $\#P$  is contained in  $\#BQP$ , as the latter includes classical verifiers. It remains to be shown that any  $\#BQP$  problem can be solved by computing a  $\#P$  function. We start from a verifier operator  $\Omega$ , Eq. (1), and wish to show that the dimension of its accepting subspace, i.e., the subspace  $\mathcal{A}$  with eigenvalues  $\geq a$ , can be computed by counting satisfying inputs to some efficiently computable boolean function. Using amplification, we can ensure that  $|\dim \mathcal{A} - \text{tr } \Omega| \leq \frac{1}{4}$ , i.e., we need to compute  $\text{tr } \Omega$  to accuracy  $\frac{1}{4}$ . This can be done using a “path integral” method, which has been used previously to show containments of quantum classes in the classical classes  $PP$  and  $\#P$  (see e.g. [13]). We rewrite  $\text{tr } \Omega = \sum_{\zeta} f(\zeta)$  as a sum over products of transition probabilities along a path  $\zeta \equiv (i_0, \dots, i_N, j_1, \dots, j_N)$ , where

$$f(\zeta) = \langle i_0 |_I \langle \mathbf{0} |_A U_1^\dagger | j_1 \rangle \langle j_1 |_T U_1^\dagger \cdots U_T^\dagger | j_T \rangle \times \quad (2)$$

$$\langle i_T |_I [ |0\rangle\langle 0|_1 \otimes \mathbb{1} ] | i_T \rangle \langle i_T |_T U_T \cdots U_1 | i_0 \rangle_I | \mathbf{0} \rangle_A$$

(cf. Fig. 1). Since any such sum over an efficiently computable  $f(\zeta)$  can be determined by counting the satisfying inputs to some boolean formula (see the Supplementary Material for details), it follows that  $\Omega$  can be computed using a single query to a black box solving  $\#P$  problems.

*Summary and discussion.*—In this work, we considered two problems: Computing the density of states and computing the ground state degeneracy of a local Hamiltonian of a spin system. In order to capture the computational difficulty of these problems we introduced the quantum complexity class  $\#BQP$ , the counting version of the class  $QMA$ . We proved that this complexity class exactly captures the difficulty of our two problems, even when restricting to local Hamiltonians on 2D lattices of qubits or to 1D chains, since all these problems are complete problems for the class  $\#BQP$  [20].

We have further shown that  $\#BQP$  is no harder than its classical counterpart  $\#P$ . In particular this implies that computing the density of states is no harder for quantum Hamiltonians than it is for classical ones. While this quantum-classical equivalence might seem surprising at the Hamiltonian level, it should be noted that the classes  $\#P$  and  $PP$  quite often form natural “upper bounds” for many quantum *and* classical problems.

What about the problem of computing the density of states for fermionic systems, such as many-electron systems? On the one hand, this problem will be still in  $\#BQP$  and thus  $\#P$ , since any local fermionic Hamiltonian can be mapped via the Jordan-Wigner transform to

a (non-local) Hamiltonian on a spin system, whose energy can still be estimated efficiently by a quantum circuit [14]. On the other hand, hardness of the problem for  $\#BQP$  can be shown e.g. by using the  $\#BQP$ -hardness of  $\#LH$ , and encoding each spin using one fermion in two modes, similar to [14]. Thus, computing the density of states for fermionic systems is a  $\#BQP$ -complete problem as well.

*Acknowledgments.*—We thank S. Aaronson, D. Gottesman, D. Gross, Y. Shi, M. van den Nest, J. Watrous and S. Zhang for helpful discussions and comments. Research at PI is supported by the Government of Canada through Industry Canada and by the Province of Ontario through the Ministry of Research and Innovation. NS is supported by the Gordon and Betty Moore Foundation through Caltech’s Center for the Physics of Information and NSF Grant No. PHY-0803371.

*Note added.*—After completion of this work, we learned that Shi and Zhang [15] have independently defined  $\#BQP$  and shown its relation to  $\#P$  using the same technique.

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  - [18] Note that the connection between  $QMA$  with a unique “good witness”, such as in our  $\#BQP$  definition, and local Hamiltonians with unique ground state and a  $1/\text{poly}(n)$  gap has been shown in [6].

- [19] Formally speaking, the reduction from  $\#BQP$  to  $\#P$  is *weakly parsimonious*, i.e., for any function  $f \in \#BQP$  there exist polynomial-time computable functions  $\alpha$  and  $\beta$ , and a function  $g \in \#P$ , such that  $f = \alpha \circ g \circ \beta$ . This differs from Karp reductions where no postprocessing is allowed,  $\alpha = \text{Id}$ , but it still only requires a single call to a  $\#P$  oracle.
- [20] Note that computing the density of states to multiplicative accuracy is less difficult, see Refs. [16, 17].