

Implications of ignorance for quantum-error-correction thresholds

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Quantum error-correcting codes have a distance parameter, conveying the minimum number of single spin errors that could cause error correction to fail. However, the success thresholds of the finite per-qubit error rate that have been proven for the likes of the toric code require them to work well beyond this limit. We argue that, without the assumption of being below the distance limit, the success of error correction is not only contingent on the noise model, but what the noise model is *believed* to be. Any discrepancy must adversely affect the threshold rate, and risks invalidating existing threshold theorems. We prove that for the two-dimensional (2D) toric code, suitable thresholds still exist by utilizing a mapping to the 2D random bond Ising model.

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I. INTRODUCTION

The inherently analog nature of quantum superpositions makes the error correction of quantum systems a formidable challenge. In principle, good error-correcting codes exist [1], and have been implemented experimentally [2]. The theoretical extension to the regime of fault-tolerance [3], which requires a threshold error rate below which error correction is successful even in the presence of faulty operations, is vital to extend existing few-qubit experiments to the realms of functional quantum processors. However, the original route towards fault-tolerance, via a concatenated hierarchy of error-correcting codes, introduces such massive overheads that these schemes are impractical for the foreseeable future.

Surface codes, such as the toric code in two dimensions [4], shift the paradigm of fault-tolerance, disposing of the hierarchical structure. The corresponding reduction in systemic overheads makes them far more promising for experimental implementation. The error-correction process only requires classical processing of the results from measurements on small blocks of neighboring qubits. The resultant error-correcting and fault-tolerant thresholds are among the highest known [5], operating in a regime well beyond that predicted by the distance of the code: for a lattice of $2N^2$ qubits, a logical error can be produced by $\frac{1}{2}N$ single qubit errors, while, for large N , almost all distributions of $N^2/10$ local errors can be successfully corrected. When operating inside the distance limit of a code, no explicit knowledge of the error mechanism is required. Beyond this limit, two different physical errors may correspond to the same error syndrome. Correcting for the wrong one could lead to a logical error. Evidently, intimate knowledge of the errors is required to determine the most likely correction for a given syndrome. However, this information may not be directly available from the syndrome measurements. Previous rigorous derivations of error-correcting and fault-tolerant thresholds have assumed perfect knowledge of the noise model (in the sense that a fault-tolerant threshold is derived under the assumption that, for instance, the error rates of two error types are equal). It is therefore important to assess the impact that this ignorance might have on error correcting thresholds. A significant impact is not expected as existing algorithms such as the minimum weight perfect matching [6] function in the biased regime.

Nevertheless, our interest here is in proof rather than numerical outcomes from (possibly nonoptimal) algorithms running on finite-sized systems.

In this paper, we examine the toric code in two dimensions, subject to a local noise model. For simplicity of exposition, error correction is assumed to be implemented perfectly, while we might have imperfect knowledge of the error model. We elucidate the influence of making assumptions about the noise model, and prove that the threshold error rates are altered, but not significantly so. Of course, this setting is not physically realistic; if we can implement the syndrome measurements perfectly, it would seem reasonable that we can also determine the noise mechanism. However, the primary purpose of this paper is simply to convey that one should do this, and, furthermore, should monitor the error model to account for any drift during an experiment. Nevertheless, if that knowledge is imperfect, an error-correcting threshold still exists. Moreover, the results presented here can, in principle, be extended to a discussion of fault-tolerance,¹ at which point we cannot know the error model perfectly.

The main technical tool that we use is the previously established connection between syndrome measurements on a noisy toric code and the random bond Ising model (RBIM) in two dimensions [5] (or three dimensions if the measurements are noisy). The phase transition of the RBIM locates the critical threshold of the toric code. In [5], this connection was established for a noise model parametrized only by p , the error rate of a known model. We extend this to a noise model that contains two error rates \tilde{p}_X and \tilde{p}_Z , and our assumptions about what these values are p_X and p_Z . The critical region of the RBIM can be determined by an ansatz [7–10] and improved upon by a renormalization style expansion [11]. The values resulting from this ansatz are numerically verified via explicit simulation of a correction algorithm, minimum weight perfect matching.

Other authors [12,13] have recently concerned themselves with the idea that two different error types, X and Z , could occur at different rates. The standard version of the toric

¹The numerical computation of useful bounds becomes a more formidable challenge, although the theory all readily extends.

code in two dimensions does not tolerate these well, with a threshold of the form $\max(p_X, p_Z) \leq p_C$, and so they have studied how one might alter the lattice geometry to better tolerate asymmetries between the parameters p_X and p_Z . The model that we choose to study here is a minor transformation of the standard two-dimensional (2D) toric code, making it more akin to Wen's model [14]. This has superior symmetry properties, vastly increasing the range of parameters for which error correction is possible. These thresholds come close to, or even exceed the quantum Hamming bound, which limits the performance of nondegenerate codes by measuring the information content of typical error sequences. While it is known that degenerate codes such as the toric code can surpass this bound [15], few instances are known.

Before we begin in earnest, let us present an initial statement that justifies our assertions on the existence of good error-correcting thresholds. In the presence of X and Z errors, if the two types of error occur independently with the same probabilities p , it was shown in [5] that error correction is possible if $p < p_C$, some threshold. In that proof, the X and Z errors are treated independently of one another. As such, if the two error rates exhibit some bias, it is clear that error correction is still possible if

$$\max(p_X, p_Z) < p_C.$$

So, this immediately shows that a threshold remains if $p_X \neq p_Z$. Moreover, any error correction algorithm that exhibits its own threshold $p < p_C^A$ and operates independently on the two error types must display a similar relationship, $\max(p_X, p_Z) < p_C^A$. This includes, for instance, the minimum weight perfect matching algorithm, for which p_C^A is very close to p_C , and it achieves this without any knowledge of what the bias is. Hence, perfect knowledge of the underlying noise model is unnecessary. However, that threshold is quite weak: Fig. 1 compares the small error-correcting region (dashed lines) that can be achieved in this way to the quantum Hamming bound (solid black line), which gives a good estimate for how well we might hope to be able to perform. The important issues are how badly the error correcting thresholds are affected by a lack of knowledge of the error model, and to what extent partial knowledge of the error parameters can benefit the threshold.

Toric code

The toric code [4] is the quintessential example of a surface code. We consider here, as in [16], a rotated version akin to the Wen code [14]. This was also studied by the authors of [17]. To define it, start from an $N \times N$ square lattice with periodic boundary conditions. Later, it will be convenient for us to term this the *primal lattice*. The dual lattice is identical, but shifted by half a lattice vector both horizontally and vertically. On the primal lattice, place a qubit in the middle of each edge. Each vertex v and face f has four neighboring qubits, two on horizontal edges E_H and two on vertical edges E_V . The measurement operators of the code are defined for each vertex and face as

$$K_v = \prod_{e \in E_H} Z_e \prod_{e \in E_V} X_e, \quad K_f = \prod_{e \in E_H} X_e \prod_{e \in E_V} Z_e.$$

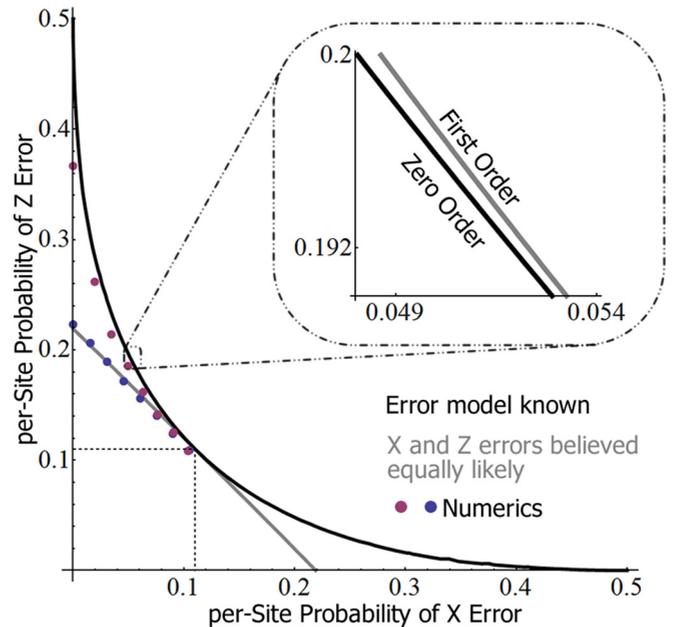


FIG. 1. (Color online) Comparison of error-correction threshold when the error model is known ($\tilde{p}_X = p_X, \tilde{p}_Z = p_Z$) and when errors are assumed equally likely $p_X = p_Z$. Plotted for comparison are numerical thresholds from minimum weight perfect matching simulations, see Sec. III. The enlarged region demonstrates the difference between the zero-order approximation of Eq. (4), and the first-order correction of Eq. (5). The dashed region is the error-correcting region for the toric code in its unrotated form.

All the terms, known as stabilizers, mutually commute and have eigenvalues ± 1 . The space of the toric code states $|\Psi_{ij}\rangle$ for $i, j \in \{0, 1\}$ are defined by the relations $K_v |\Psi_{ij}\rangle = |\Psi_{ij}\rangle$ and $K_f |\Psi_{ij}\rangle = |\Psi_{ij}\rangle$ for all v, f . There are $2N^2$ qubits and $2N^2 - 2$ independent stabilizers, leaving a four-fold degeneracy (the indices i, j) that represents two logical qubits. The two logical Pauli Z^L (X^L) operators correspond to products of Z (X) operations along a single column (row), looping around the entire torus. There are two inequivalent columns (rows), composed of either horizontal or vertical edges. Starting from a logical state $|\Psi_{ij}\rangle$, and applying continuous segments of X and Z operators, it is possible to form closed loops (meaning that all stabilizers return +1 expectation). Provided those loops are topologically trivial (i.e., they do not form loops around the torus), the state is the same as the original one while nontrivial loops correspond to logical errors. This degeneracy of the code means that if a large set of errors has arisen, it is not necessary to establish exactly which errors occurred to correct for them; one only has to form the closed loops which are most likely to be trivial.

We consider an error model of X and Z errors acting independently on each site, with probabilities \tilde{p}_X and \tilde{p}_Z , respectively: a single qubit state ρ undergoes

$$\rho \mapsto \mathcal{E}_Z(\mathcal{E}_X(\rho)); \quad \mathcal{E}_\sigma(\rho) = (1 - \tilde{p}_\sigma)\rho + \tilde{p}_\sigma \sigma \rho \sigma.$$

We will give equivalent results for the extended model

$$\mathcal{E}(\rho) = (1 - \tilde{q}_X - \tilde{q}_Z - \tilde{q}_Y)\rho + \tilde{q}_X X \rho X + \tilde{q}_Z Z \rho Z + \tilde{q}_Y Y \rho Y$$

in Sec. IV. The restricted error model of X and Z errors only means that there are two independent sets of errors (anyons) that can never interact; those detected by the $\{K_f\}$ and $\{K_v\}$, respectively. By symmetry, it suffices to consider just one of these sets, say $\{K_v\}$. To see this, let us denote by L the lattice on which the toric code is defined (with the qubits in the middles of the edges). Let L_1 be a copy of L and L_2 be the dual of L_1 (i.e., in the case of the periodic square lattice, the same lattice but shifted both horizontally and vertically by half a unit). With the edges q (corresponding to a qubit on L) of each of the lattices L_i we associate a variable $\tau_q^i \in \pm 1$ in the following way:

Qubit type	Error type	Assignment
$q \in V$	X	$\tau_q^2 = -1$
$q \in V$	Z	$\tau_q^1 = -1$
$q \in H$	X	$\tau_q^1 = -1$
$q \in H$	Z	$\tau_q^2 = -1$.

Note that X and Z errors affect the two lattices equally.² All other variables τ are set to 1. Note that this means that X errors are specified by horizontal edges of L_1 and L_2 , while Z errors are specified by vertical edges.

II. RANDOM BOND ISING MODEL

In [5], a connection was proven between the ability to correct errors arising on the toric code and the existence of a phase transition in the random bond Ising model (RBIM). This was done for both perfect (the 2D RBIM) and imperfect stabilizer measurement [three-dimensional (3D) RBIM], assuming that X and Z errors each occur with probability p independently on each lattice site, and assuming that p is known. For simplicity, we will only consider the case of perfect stabilizer measurement, our aim being to relax the assumptions on the knowledge of the error rates, and their equality. To emphasize the difference with the actual error rates (\tilde{p}_X and \tilde{p}_Z), the assumed error rates are denoted by p_X and p_Z , respectively. One should only be able to achieve the optimal recovery specified by [5] if the nature of the noise is known exactly.

Whether (or not) the error rates \tilde{p}_X and \tilde{p}_Z might be inferred from the syndrome measurements, it is clearly feasible to monitor the error rates to get a good estimate. We do not claim that they should be completely unknown. Rather, our purpose here is twofold: (i) find the error correcting threshold when the rates $p_X = \tilde{p}_X$ and $p_Z = \tilde{p}_Z$ and (ii) convey that it is important to be working as close to the conditions $p_X = \tilde{p}_X$ and $p_Z = \tilde{p}_Z$ as possible but that, nevertheless, an error correcting threshold still exists, i.e., that it is sufficient to have an estimate on the error rates, rather than needing an exact characterization of the full error model.

For one set of errors, say those affecting the $\{K_v\}$, once the error syndrome has been extracted by measuring the

stabilizers, the aim of error correction is to apply a set of operations that reset all the stabilizers to $+1$. Relative to the state that was initially encoded, there are only four inequivalent consequences ($\mathbb{1}$, X_1^L , Z_2^L , and $X_1^L Z_2^L$) of the correction. Four corresponding corrections, error strings E_i , $i = 0, \dots, 3$, can be identified and need to be assigned a likelihood of having arisen

$$p_i = \sum_{C \in S} p(E_i \cup C)$$

according to the assumptions on the noise model, where the set S corresponds to all trivial loops. If error correction is possible, then the expectation of the probability of getting the right answer over all actual error configurations should tend to 1 in the limit of large system size N , while the other probabilities should vanish. When error correction fails, all the p_i will be similar.

Let τ^0 be a set of variables ± 1 for each qubit, corresponding to whether or not a rotation is applied in the correction E_0 . Similarly, τ_C^0 is the set due to $E_0 \cup C$. The set of all closed loops is conveniently described by introducing variables $\sigma_i \in \{\pm 1\}$ for each vertex of the dual lattice [5, 18]. A qubit q on an edge of the primal lattice has two neighboring vertices of the dual lattice v_1^q and v_2^q : $\tau_{C,q}^0 = \tau_q^0 \sigma_{v_1^q} \sigma_{v_2^q}$, where it now suffices to sum over the variables σ_i without restriction. For probabilities

$$p_q = \begin{cases} p_Z & q \in V, \\ p_X & q \in H, \end{cases}$$

we assign the probability of a given error string as

$$p(\tau_C^0) = \prod_q (1 - p_q)^{(1 + \tau_{C,q}^0)/2} p_q^{(1 - \tau_{C,q}^0)/2}.$$

Removing a common factor, we have

$$\prod_q \left(\frac{1 - p_q}{p_q} \right)^{\tau_{C,q}^0/2}.$$

By defining

$$\frac{1 - p_X}{p_X} = e^{2J_H} \quad \frac{1 - p_Z}{p_Z} = e^{2J_V},$$

the probability p_0 is proportional to

$$Z_0 = \sum_{\vec{\sigma}} e^{H(\vec{\sigma})}$$

with

$$H(\vec{\sigma}) = \sum_{q \in H} (\tau_q^0 J_H) \sigma_{v_1^q} \sigma_{v_2^q} + \sum_{q \in V} (\tau_q^0 J_V) \sigma_{v_1^q} \sigma_{v_2^q}. \quad (1)$$

This is the Hamiltonian of the $\pm J$ random bond Ising model on a square lattice, where the vector τ^0 arises from the actual errors that occurred, and the coupling strengths constitute our assignment of the likelihood of different configurations. The transition in behavior of the probabilities between successful correction (in asymptotically all instances of the syndromes) and failure corresponds to a discontinuity of the free energy $F = \ln Z$ of this model, where $Z = \sum_i Z_i$. This is the well-studied phase transition in the 2D RBIM. To determine the phase transitions, we first establish the duality of the nonrandom version of the Ising model, and subsequently extend it via

²In the usual (unrotated) toric code, all X errors give the values of τ_q^1 , and all Z errors give the values of τ_q^2 .

the replica method to account for the configurational average in the random version.

A. Duality

Our analysis starts by considering the nonrandom bond Ising model, aiming to reproduce the result that the model has a self-dual point [18]. While true for all configurations τ^0 , we set $\tau_q^0 = 1 \forall q$ for simplicity. Why is duality interesting? It means that as a single parameter p is varied, there is a certain symmetry present such that the model behaves the same also at another point $f(p)$. Assuming the model has exactly one phase transition, the only place that this can occur is the self-dual point $p = f(p)$ [if it occurred at any other point p , it would also occur at point $f(p)$]. We start by expressing

$$Z_0 = \sum_{\vec{\sigma}} \prod_q u_q (\sigma_{v_1^q} \sigma_{v_2^q}).$$

This is just writing exactly what we had before but with

$$u_q(\tau) = \begin{cases} e^{J_H \tau} & q \in H, \\ e^{J_V \tau} & q \in V. \end{cases}$$

Now, if we define

$$\sqrt{2}u_q^*(s) = \begin{cases} e^{J_H s} + (-1)^s e^{-J_H} & q \in H, \\ e^{J_V s} + (-1)^s e^{-J_V} & q \in V, \end{cases}$$

for $s \in \{0, 1\}$, then we can express

$$\sqrt{2}u_q(\tau) = \sum_{s_q \in \{0, 1\}} u_q^*(s_q) (-1)^{s_q (\varepsilon_{v_1^q} + \varepsilon_{v_2^q})},$$

having replaced σ_i with $2\varepsilon_i - 1$ such that $\varepsilon_i \in \{0, 1\}$. Now consider what happens when we examine the product of u_q , expressed as above, for all edges of L leading out of a given vertex i , and performing the sum in Z_0 over that particular ε_i , i.e.,

$$Z_0 = \sum_{\vec{\sigma}} \prod_{v_i^r} u_r \sum_{\sigma_i \in \{\pm 1\}} \prod_{q: v_1^q = i} u_q$$

has that

$$\begin{aligned} & \sum_{\sigma_i} \prod_{q: v_1^q = i} u_q (\sigma_i \sigma_{v_2^q}) \\ &= \frac{1}{4} \sum_{\varepsilon_i \in \{0, 1\}} \prod_{q: v_1^q = i} \sum_{s_q \in \{0, 1\}} u_q^*(s_q) (-1)^{s_q (\varepsilon_i + \varepsilon_{v_2^q})}, \end{aligned}$$

and contained within this is

$$\sum_{\varepsilon_i \in \{0, 1\}} (-1)^{\varepsilon_i \sum_j s_j},$$

which gives a value 2 if

$$\sum_j s_j \bmod 2 = 0,$$

and 0 otherwise. The sum j is over the vertices such that we include all q with vertices i and j . This is entirely equivalent to the product of s_j having to be +1 around vertices or, in other words, around closed loops of the dual lattice. Hence,

$$Z_0 = \sum_s \prod_q u_q^*,$$

where $'$ denotes the restricted sum only for satisfying assignments around closed loops of the dual lattice.³ This is exactly the same as on the primal lattice, where we could have written

$$Z_0 = \sum_{\tau} \prod_q u_q$$

with $\tau_q = \sigma_{v_1^q} \sigma_{v_2^q}$, and $'$ indicating a restricted sum only for satisfying assignments around closed loops of the primal lattice. Thus, the model is self-dual when

$$u_q(\tau) = u_q^*[(1 - \tau)/2] \quad (2)$$

for all q and all $\tau \in \{\pm 1\}$, remembering that a horizontal edge on the primal lattice corresponds to a vertical edge on the dual lattice. It turns out that the only condition for the self-dual point is

$$e^{-2J_H} = \tanh(J_V),$$

satisfying all four equations (2) simultaneously.

B. Replica method

The above duality was proven without any randomness present. To deal with the randomness of the bonds, we must take a configurational average over the possible values of τ^0 , and use it to determine any discontinuity in the free energy $\ln Z$. The way that we approach calculating this is to consider n parallel copies of the model (all with the same configuration of ± 1 bonds). The partition function of all n parallel copies is just Z^n , which is readily calculated for positive integers n . In the limit $n \rightarrow 0$ exists, then

$$\ln Z = \left\langle \lim_{n \rightarrow 0} \frac{Z^n - 1}{n} \right\rangle.$$

For n copies, we can perform the same duality studies as we did above for a single copy. However, the function $u_q(\tau)$ with $\tau \in \{\pm 1\}$ must be replaced with $u_q(\tau)$ with $\tau \in \{\pm 1\}^n$, i.e., there is a value of ± 1 for the bond q in each copy. The same happens for u^* , and self-duality only arises if

$$u_q(\tau) = u_q^*[(1 - \tau)/2]$$

for all $\tau \in \{\pm 1\}^n$ simultaneously. This problem reduces slightly because these expressions only depend on the number of the number of +1's in the vector τ , so we only need to test equality for $n + 1$ cases rather than 2^n . Let us take a vector \vec{x} of $n + 1$ elements, and assign to element x_p the value of $u_q(\tau)$ when τ contains $p - 1$ values.

$$x_k^H = \tilde{p}_X e^{(n-2k)J_H} + (1 - \tilde{p}_X) e^{-(n-2k)J_H},$$

$$x_k^{*H} = \sqrt{2^n} \cosh^n(J_H) \tanh^k(-J_H) \{1 + [1 + (-1)^k] \tilde{p}_X\}.$$

The V versions are equivalent, with $\tilde{p}_X \mapsto \tilde{p}_Z$ and $J_H \mapsto J_V$. Unfortunately, after averaging over the possible random bond assignments τ_q^0 , there are no values of J_H and J_V such that

³The factors of 2 conveniently cancel: for an $n \times m$ lattice, there are $2nm$ edges, and hence a factor $\sqrt{2^{-2nm}}$ appears when replacing the u with u^* for each edge. The compensating factor is the factor of 2 that arises for each vertex (of which there are nm) from the sum over ε_i .

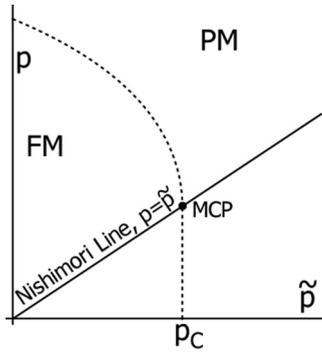


FIG. 2. Phase diagram of the 2D RBIM. The dashed line separates the ferromagnetic (FM) and paramagnetic (PM) phases. Where the Nishimori line intersects this is the multicritical point.

$\bar{x} = \bar{x}^*$. Had we been able to, then this is the limit we would have taken as $n \rightarrow \infty$ to find the phase transition. Instead, we follow the approach of [7–10], in which it was conjectured that the critical point is approximated by

$$x_0^H x_0^V = x_0^{*H} x_0^{*V}. \quad (3)$$

Taking the limit $n \rightarrow 0$ yields

$$\begin{aligned} \tilde{p}_H \log_2 p_H + (1 - \tilde{p}_H) \log_2(1 - p_H) \\ + \tilde{p}_V \log_2 p_V + (1 - \tilde{p}_V) \log_2(1 - p_V) = -1. \end{aligned} \quad (4)$$

The ansatz of Eq. (3) was postulated in [7–10] specifically to work at the multicritical points of the RBIM ($\tilde{p}_X = p_X$ and $\tilde{p}_Z = p_Z$), and this was justified by the existence of various symmetries. It has subsequently been numerically tested extensively within this regime, and the asymmetric case of $p_X \neq p_Z$ [10]. In this case, Eq. (4) coincides exactly with the quantum Hamming bound. Moreover, the specific instance of $\tilde{p}_X = \tilde{p}_Z = p_X = p_Z$ transforms to the Nishimori line of the RBIM, and reveals the critical probability quoted in [5] via correspondence to the multicritical point of the RBIM.

Practically, we will only ever be able to estimate the parameters p_X and p_Z , rather than determine them exactly. Consider the worst possible case, in which we determine the frequency p of stabilizers being -1 , and assume that $p_X = p_Z$. How detrimental is this to the threshold? Equation (4) reduces to an effective homogeneous system with $2\tilde{p} = \tilde{p}_X + \tilde{p}_Z$ and $(1 - 2p)^2 = (1 - 2\tilde{p}_X)(1 - 2\tilde{p}_Z)$. This point lies on or above the Nishimori line ($\tilde{p} > p$, see Fig. 2). While Eq. (3) was not originally proposed to function in this regime, detailed studies [11] confirmed that in the homogeneous case, above the Nishimori line, the approximation is a good one. Hence, our analysis remains reliable. Alternatively, having reduced to the homogeneous case, the critical probability cannot be larger than that at the multicritical point $p_C \approx 0.1092$. Hence, setting $p = \tilde{p}$ yields $\tilde{p}_X + \tilde{p}_Z < 2p_C$, as compared to the nontransformed version which only successfully corrects if $\max(\tilde{p}_X, \tilde{p}_Z) < p_C$. The transformed version has more natural symmetry properties and negates the requirement of recent studies [12,13] to adjust the lattice geometry for each different asymmetry between \tilde{p}_X and \tilde{p}_Z . Figure 1 shows that for all parameter values, a finite per-qubit error rate threshold remains and is superior to our original crude estimate (the dashed lines). Provided a sufficiently accurate estimate of the parameters

of the error model is made, the error threshold is essentially unaffected.

Although the conjecture of Eq. (3) compares favorably with numerical estimates on the square lattice, there is known to be a discrepancy with some exact renormalization group calculations on hierarchical lattices [19]. To account for this, a renormalization inspired expansion was introduced in [11,20] to account for corrections. While most natural for hierarchical lattices [20], it has been extended to square lattices and achieves an even tighter match with previous numerical results using only a first-order correction [11]. This technique proceeds by replacing the term e^{J_H} in x_0^H (and, similarly, the term $e^{J_V} + e^{-J_H}$ in x_0^{*H}) with an equivalent effective weight arising from an averaging effect over several neighboring spins. The idea is that a differing order of correction can be calculated by considering larger and larger neighborhoods. In the large neighborhood limit, one must certainly recover the true behavior of the model. The first-order approximation considers only nearest neighbors, locating the critical point at

$$\begin{aligned} 2 = \frac{1}{2} \sum_{\eta=\pm 1} [1 + \eta(1 - 2p)^4] \log_2 [1 + \eta(1 - 2p)^4] \\ - \sum_{n=0}^1 \sum_{m=0}^2 \binom{2}{m} a_{nm}(\tilde{p}_X, \tilde{p}_Z) \log_2 a_{nm}(p_X, p_Z), \end{aligned} \quad (5)$$

where

$$a_{nm}(r, s) = \frac{r^n s^m}{(1 - r)^{n-2} (1 - s)^{m-2}} + \frac{(1 - r)^n (1 - s)^m}{r^{n-2} s^{m-2}}.$$

This is a generalized version of Eq. (45) in [11], accounting for an asymmetry between X and Z errors. The threshold values are barely impacted, although Fig. 1 indicates that they can surpass the zero-order approximation. This violation increases at second order, suggesting that this is not a finite-sized computational effect, and that this code does indeed have an error correcting threshold (in certain regimes) which exceeds the quantum Hamming bound. This is by no means forbidden (the quantum Hamming bound only applies to nondegenerate error-correcting codes, while the toric code is degenerate), but few examples are known [15].

III. MINIMUM WEIGHT PERFECT MATCHING

Given the near-vertical phase boundary of the RBIM below the Nishimori line (Fig. 2), any error correction strategy which assumes a lower “temperature” has an almost identical critical probability. In particular, the zero temperature case corresponds to correcting by minimum weight perfect matching;⁴ an efficient algorithm which is readily implemented. As such, it provides a lower bound on the threshold fidelities for verification of the previous results. For this purpose, the (im)practicalities of its application [21] are irrelevant.

⁴In this context, temperature is a mathematical parameter of the mapping to the RBIM and has no physical analog. Zero temperature corresponds to the limit $J \rightarrow \infty$, meaning the only term in Z_0 that is worth considering is the $\vec{\sigma}$ that gives the smallest value.

The input to a minimum weight perfect matching algorithm is a set of vertices. In the present case, these vertices correspond to the stabilizers which give -1 values (i.e., the locations of anyons). We must then assign a weight for every pairing of two anyons. We will justify a weighting function momentarily. The algorithm then outputs the way in which the anyons can be paired up (i.e., how they might annihilate each other) such that the total weight is minimized. The idea is to make this correction correspond to the most probable set of operations that could have created that distribution of anyons. If a given anyon pair are separated by l_H and l_V in the horizontal and vertical directions, then, as a minimum, they must have been created by l_H X errors and l_V Z errors. Hence, we assign a minimum probability of

$$\left(\frac{p_X}{1-p_X}\right)^{l_H} \left(\frac{p_Z}{1-p_Z}\right)^{l_V}$$

to that combination. So, if we take a particular way of pairing up all the anyons, the probability that such a combination arose was the product of all the individual pair-wise probabilities. We want to find the combination that minimizes that product, but that is the same as finding the combination that minimizes the sum of corresponding logarithms

$$l_H \ln \left(\frac{p_X}{1-p_X}\right) + l_V \ln \left(\frac{p_Z}{1-p_Z}\right).$$

Hence these constitute the weights that we must minimize the total of.

Our simulation, with the results depicted in Fig. 1, functions by considering an $N \times N$ lattice where $N = 100$. We implemented an error model that created X and Z errors on each qubit with probabilities \tilde{p}_X and \tilde{p}_Z , respectively. Having ascertained the positions of each error (i.e., which stabilizers anticommute with the errors), we assigned weights between a pair of vertices separated by l_H and l_V in the horizontal and vertical directions as

$$l_H \ln \left(\frac{p_X}{1-p_X}\right) + l_V \ln \left(\frac{p_Z}{1-p_Z}\right).$$

The BLOSSOM V algorithm [22] was then used to perform the minimum weight perfect matching. For a fixed ratio \tilde{p}_X/\tilde{p}_Z , the fraction of 500 different realizations of an error distribution giving a logical error was computed for varying error rates, enabling the determination of the failure probability (the threshold at which a transition in logical error rate from 0 to 50% occurs). Similar numerics, for a perfectly identified error model, are present in [12].

IV. GENERALIZED MODEL

The previous analysis can be repeated for a more general error model of

$$\mathcal{E}(\rho) = (1 - \tilde{q}_X - \tilde{q}_Y - \tilde{q}_Z)\rho + \tilde{q}_X X\rho X + \tilde{q}_Y Y\rho Y + \tilde{q}_Z Z\rho Z.$$

Note that we use q to distinguish from the previous p . The previous distribution had

$$q_X = p_X(1 - p_Z), \quad q_Z = p_Z(1 - p_X), \quad q_Y = p_X p_Z.$$

With Y errors present, it is not possible to divide the original system into two independent systems, as we did before. Nevertheless, we can still express the probability of successful error correction as being related to the phase transition of the model

$$Z_0 = \sum_{\sigma} e^H$$

with

$$H = \sum_{q \in H} \tau_q^1 J_H \sigma_i \sigma_j + \tau_q^2 J_V \sigma_{i'} \sigma_{j'} + \tau_q^1 \tau_q^2 J_Y \sigma_i \sigma_j \sigma_{i'} \sigma_{j'},$$

where

$$\begin{aligned} e^{AJ_H} &= \frac{(1 - q_X - q_Y - q_Z)q_Z}{q_Y q_X}, \\ e^{AJ_V} &= \frac{(1 - q_X - q_Y - q_Z)q_X}{q_Y q_Z}, \\ e^{AJ_Y} &= \frac{(1 - q_X - q_Y - q_Z)q_Y}{q_X q_Z}. \end{aligned}$$

The τ_q^i have a sign distribution specified by the error model. This is exactly the model derived in [23]. The duality and replica arguments follow in much the same way. For instance, the duality transformation of a single copy is described by

$$\underline{u} = (e^{J_H + J_V + J_Y}, e^{J_H - J_V - J_Y}, e^{-J_H + J_V - J_Y}, e^{-J_H - J_V + J_Y}),$$

$$\underline{u}^* = \frac{1}{2} \begin{pmatrix} 1 & 1 & 1 & 1 \\ 1 & -1 & 1 & -1 \\ 1 & 1 & -1 & -1 \\ 1 & -1 & -1 & 1 \end{pmatrix} \underline{u},$$

and equality (i.e., self-duality) can be generated if, for example, $J_H = J_V$ and

$$e^{-2J_Y} = \sinh(2J_H).$$

Finally, we get the equivalent of Eq. (4), i.e., the zero-order approximation, for this generalized model

$$\begin{aligned} (1 - \tilde{q}_X - \tilde{q}_Y - \tilde{q}_Z) \log_2(1 - \tilde{q}_X - \tilde{q}_Y - \tilde{q}_Z) \\ + \sum_{\sigma \in \{X, Y, Z\}} \tilde{q}_\sigma \log_2 \tilde{q}_\sigma = -1. \end{aligned}$$

One application is assessing how well minimum weight perfect matching might perform on depolarizing noise, as compared to the optimal. If $q_X = q_Y = q_Z = \tilde{q}_X = \tilde{q}_Y = \tilde{q}_Z$, i.e., depolarizing noise that we have perfectly identified, one has to solve the equation

$$(1 - 3q) \log_2(1 - 3q) + 3q \log_2(q) = -1$$

to find the critical $3q = 0.18929$, which replicates the value given in [23]. This is the best that error correction could achieve. What about minimum weight perfect matching? Since it is not capable of taking the correlations introduced by Y into account, it is not expected to be tight with the optimal correction. So, it effectively proceeds by making the assumption that $q_Y = p^2$, $q_X = q_Z = p(1 - p)$, $\tilde{q}_X = \tilde{q}_Y = \tilde{q}_Z = \tilde{q}$. Moreover, we know that the performance of the algorithm is very similar to that of the critical point of the model parametrized in this way. So, it suffices to solve for

the critical point again, which is described by

$$-\frac{1}{2} = (1 - 2q) \log_2(1 - p) + 2q \log_2(p),$$

and the largest value of q is given by setting $p = 2q$, which reveals that $3q = 0.165$. This compares favorably with previous numerical estimates [24].

V. CONCLUSION

This paper has investigated how any discrepancy between the actual and assumed noise models can influence the error-correcting threshold of a code if operating in a regime beyond that specified by the code's distance. The principle is universal, although it has been demonstrated for a specific class of local

noise acting on the toric code in two dimensions. It is important to note that the rotated version of the toric code examined here exhibits a significant enhancement in robustness with respect to asymmetries between X and Z error rates, and can even surpass the quantum Hamming bound in some regimes. The concern that the threshold value might be adversely affected, as raised in [16], is unfounded; good approximations to the error parameters of a system lead to a negligible change in the threshold error rates.

We have also detailed how correlations due to Y errors can be incorporated into the analysis [23]. In principle, the present results can be generalized to the situation of noisy measurements [5,9], adding a third dimension to model to the RBIM, to derive fault-tolerant thresholds. It is anticipated that the conclusions of this paper should follow similarly.

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- [1] A. R. Calderbank and P. W. Shor, *Phys. Rev. A* **54**, 1098 (1996).
 - [2] J. Zhang, R. Laflamme, and D. Suter, *Phys. Rev. Lett.* **109**, 100503 (2012); P. Schindler, J. T. Barreiro, T. Monz, V. Nebendahl, D. Nigg, M. Chwalla, M. Hennrich, and R. Blatt, *Science* **332**, 1059 (2011).
 - [3] P. Aliferis, D. Gottesman, and J. Preskill, *Quantum Inf. Comput.* **6**, 97 (2006).
 - [4] A. Kitaev, *Ann. Phys. (NY)* **321**, 2 (2006).
 - [5] E. Dennis, A. Kitaev, A. Landahl, and J. Preskill, *J. Math. Phys.* **43**, 4452 (2002).
 - [6] C. Wang, J. Harrington, and J. Preskill, *Ann. Phys. (NY)* **303**, 31 (2003); R. Raussendorf and J. Harrington, *Phys. Rev. Lett.* **98**, 190504 (2007); D. S. Wang, A. G. Fowler, A. M. Stephens, and L. C. L. Hollenberg, *Quantum Inf. Comput.* **10**, 456 (2010).
 - [7] H. Nishimori and K. Nemoto, *J. Phys. Soc. Jpn.* **71**, 1198 (2002).
 - [8] J.-M. Maillard, K. Nemoto, and H. Nishimori, *J. Phys. A: Math. Gen.* **36**, 9799 (2003).
 - [9] K. Takeda and H. Nishimori, *Nucl. Phys. B* **686**, 377 (2004).
 - [10] K. Takeda, T. Sasamoto, and H. Nishimori, *J. Phys. A: Math. Gen.* **38**, 3751 (2005).
 - [11] M. Ohzeki, *Phys. Rev. E* **79**, 021129 (2009).
 - [12] B. Röthlisberger, J. R. Wootton, R. M. Heath, J. K. Pachos, and D. Loss, *Phys. Rev. A* **85**, 022313 (2012).
 - [13] K. Fujii and Y. Tokunaga, *Phys. Rev. A* **86**, 020303 (2012).
 - [14] X.-G. Wen, *Phys. Rev. Lett.* **90**, 016803 (2003).
 - [15] D. P. DiVincenzo, P. W. Shor, and J. A. Smolin, *Phys. Rev. A* **57**, 830 (1998).
 - [16] A. Kay, *Phys. Rev. Lett.* **107**, 270502 (2011).
 - [17] H. Bombin, *Phys. Rev. Lett.* **105**, 030403 (2010).
 - [18] F. Y. Wu and Y. K. Wang, *J. Math. Phys.* **17**, 439 (1976).
 - [19] M. Hinczewski and A. N. Berker, *Phys. Rev. B* **72**, 144402 (2005).
 - [20] M. Ohzeki, H. Nishimori, and A. N. Berker, *Phys. Rev. E* **77**, 061116 (2008).
 - [21] G. Duclos-Cianci and D. Poulin, *Information Theory Workshop, IEEE*, 1 (2010), [arXiv:1006.1362](https://arxiv.org/abs/1006.1362)[quant-ph].
 - [22] V. Kolmogorov, *Math. Prog. Comp.* **1**, 43 (2009).
 - [23] H. Bombin, R. S. Andrist, M. Ohzeki, H. G. Katzgraber, and M. A. Martin-Delgado, *Phys. Rev. X* **2**, 021004 (2012).
 - [24] G. Duclos-Cianci and D. Poulin, *Phys. Rev. Lett.* **104**, 050504 (2010).