Two-dimensional Potts antiferromagnets with a phase transition at arbitrarily large $q$

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(Received 23 October 2012; published 24 January 2013)

We exhibit infinite families of two-dimensional lattices (some of which are triangulations or quadrangulations of the plane) on which the $q$-state Potts antiferromagnet has a finite-temperature phase transition at arbitrarily large values of $q$. This unexpected result is proven rigorously by using a Peierls argument to measure the entropic advantage of sublattice long-range order. Additional numerical data are obtained using transfer matrices, Monte Carlo simulation, and a high-precision graph-theoretic method.

DOI: 10.1103/PhysRevE.87.012136 PACS number(s): 64.60.De, 05.50.+q, 11.10.Kk, 64.60.Cn

I. INTRODUCTION

The $q$-state Potts model [1,2] plays an important role in the theory of critical phenomena, especially in two dimensions [3–5], and has applications to various condensed-matter systems [2]. Ferromagnetic Potts models are by now fairly well understood, thanks to universality, but the behavior of antiferromagnetic Potts models depends strongly on the microscopic lattice structure, so that many basic questions about the phase diagram and critical exponents must be investigated case-by-case. In this article we prove the unexpected existence of phase transitions for some two-dimensional (2D) $q$-state Potts antiferromagnets at arbitrarily large values of $q$.

For Potts antiferromagnets one expects that for each lattice $\mathcal{L}$ there is a value $q_c(\mathcal{L})$ [possibly noninteger] such that for $q > q_c(\mathcal{L})$ the model has an exponential decay of correlations at all temperatures including zero, while for $q = q_c(\mathcal{L})$ there is a zero-temperature critical point. The first task, for any lattice, is thus to determine $q_c$.

Some 2D antiferromagnetic models at zero temperature can be mapped exactly onto a “height” model [6,7]. Since the height model must either be in a “smooth” (ordered) or “rough” (massless) phase, the corresponding zero-temperature spin model must either be ordered or critical, never disordered. Until now it has seemed that the most common case is criticality [8].

In particular, when the $q$-state zero-temperature Potts antiferromagnet (AF) on a 2D lattice $\mathcal{L}$ admits a height representation, one ordinarily expects that $q = q_c(\mathcal{L})$. This prediction is confirmed in most heretofore-studied cases: three-state square-lattice [6,9–11], three-state kagome [12,13], four-state triangular [14], and four-state on the line graph of the square lattice [13,15]. Until recently the only known exception was the triangular Ising AF [16].

Kotecký, Salas, and Sokal (KSS) [17] observed that the height mapping employed for the three-state Potts AF on the square lattice [6] carries over unchanged to any plane quadrangulation and Moore and Newman [14] observed that the height mapping employed for the four-state Potts AF on the triangular lattice carries over unchanged to any Eulerian plane triangulation (a graph is called Eulerian if all vertices have even degree). One therefore expects naively that $q_c = 3$ for every (periodic) plane quadrangulation, and that $q_c = 4$ for every (periodic) Eulerian plane triangulation.

Surprisingly, these predictions are false. KSS [17] proved rigorously that the three-state AF on the diced lattice (which is a quadrangulation) has a phase transition at finite temperature (see also [18]); numerical estimates from transfer matrices yield $q_c(\text{diced}) \approx 3.45$ [19]. Likewise, we recently [20] provided analytic arguments (falling short, however, of a rigorous proof) that on any Eulerian plane triangulation in which one sublattice consists entirely of vertices of degree 4, the four-state AF has a phase transition at finite temperature, so that $q_c > 4$. We also presented transfer-matrix and Monte Carlo data confirming these predictions for the union-jack (UJ) and bisected hexagonal (BH) lattices, leading to the estimates $q_c(\text{UJ}) \approx 4.33$ and $q_c(\text{BH}) \approx 5.40$.
These results suggest the obvious question: How large can \( q_c \) be on a plane quadrangulation (or Eulerian plane triangulation)? The answers are clearly larger than 3 or 4, respectively; but how much larger?

In this article we shall give a rigorous proof of the unexpected answer: We exhibit infinite classes of plane quadrangulations and Eulerian plane triangulations on which \( q_c \) can take arbitrarily large values. We shall also complement this rigorous proof with detailed quantitative data from transfer matrices, Monte Carlo simulations, and a powerful graph-theoretic approach developed recently by Jacobsen and Scullard [21].

The models studied here provide new examples of entropically driven long-range order [17,18,20,22]: The ferromagnetic ordering of spins on one sublattice is favored because it increases the freedom of choice of spins on the other sublattice(s). But though this idea is intuitively appealing, it is usually difficult to determine quantitatively, in any specific case, whether the entropic penalty for interfaces between domains of differently ordered spins on the first sublattice is large enough to produce long-range order. Moreover, one expects that this penalty decreases with increasing \( q \). In the examples given here, by contrast, we are able to prove that the penalty can be made arbitrarily strong and hence operative at arbitrarily large \( q \).

II. LATTICES \( G_n \) AND \( H_n \)

Let \( G_n \) be obtained from the square (SQ) lattice by replacing each edge with \( n \) two-edge paths in parallel and let \( H_n \) be obtained from \( G_n \) by connecting each group of \( n \) "new" vertices with an \((n-1)\)-edge path (see Fig. 1). Resumming over the spins on the "new" vertices [23], it is easy to show that the \( q \)-state Potts model on \( G_n \) or \( H_n \) with nearest-neighbor coupling \( v = e^{f} - 1 \) is equivalent to a SQ-lattice Potts model with a suitable coupling \( v_{\text{eff}}(q,v) \) [24]; moreover, for \( q > 2 \) (respectively \( q > 3 \)) an AF model (\( -1 \leq v \leq 0 \)) on \( G_n \) (respectively \( H_n \)) maps onto a ferromagnetic model (\( v_{\text{eff}} \geq 0 \)) on the SQ lattice. Concretely, for the zero-temperature AF (\( v = -1 \)) we have

\[
\begin{align*}
    v_{\text{eff}}^{G_n}(q, -1) &= \left( \frac{q-1}{q-2} \right)^n - 1, \\
    v_{\text{eff}}^{H_n}(q, -1) &= \left( \frac{q-1}{q-2} \right)^{n-1} - 1.
\end{align*}
\]

Setting \( v_{\text{eff}} \) equal to the SQ-lattice ferromagnetic critical point \( v_c(\text{SQ}) = \sqrt{q} [3,25] \), we obtain \( q_c \) for \( G_n \) and \( H_n \); they have the large-\( n \) asymptotic behavior

\[
q_c(G_n) \approx q_c(H_n) \approx \frac{2n}{W(2n)} + O((n/\log n)^{1/2}).
\]

III. MODIFIED LATTICES

Starting from \( G_n \) or \( H_n \), insert a new vertex into each octagonal face and connect it to the four surrounding vertices of the original SQ lattice, to the four "new" vertices, or to all eight vertices; call these modifications \( ' \), \( '' \), and \( ''' \), respectively. In particular, \( G'_n \) and \( H''_n \) are quadrangulations, and \( H'''_n \) is an Eulerian triangulation (Fig. 2).

If we integrate out the spins at the vertices placed into the octagonal faces, we obtain the model on \( G_n \) or \( H_n \) perturbed by a four-spin or eight-spin interaction. When \( q \) is large, this interaction is weak (of order \( 1/q \)) because its Boltzmann weight is bounded between a maximum value of \( q \) and a minimum value of \( q-4 \) or \( q-8 \). We therefore expect that the new edges will have a negligible effect on the phase transition when \( q \) is large, and that all the modified lattices will have \( q_c(n) \) whose large-\( n \) behavior is essentially identical to Eq. (3). Let us now sketch a rigorous proof [28] of this assertion.

IV. PROOF OF PHASE TRANSITION

Recall first how one proves, using the Peierls argument, the existence of ferromagnetic long-range order (FLRO) at low temperature in the \( q \)-state Potts ferromagnet on the SQ lattice. The Peierls contours are defined as the connected components of the union of all bonds on the dual SQ lattice that...
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FIG. 4. (Color online) Estimated phase boundaries for the \( q = 4, 8, 16, 32 \) Potts models on the UJ lattice, from the Jacobsen-Scullard method (blue curve and numerical values of \( v_{\square, c} \) at \( v_x = -1 \)) and Monte Carlo simulations (black points).

FIG. 5. (Color online) Estimates of \( q_c(n) \) for the lattices \( G_n, G'_n, H_n, H'_n \) divided by their large-\( n \) asymptote \( 2n/W(2n) \).

of reflection positivity and chessboard estimates \cite{29}. Without going into details of the needed adaptations of this standard technique for our case (see \cite{28}), we mention only that the final bound on the probability of occurrence of a contour \( \gamma \) is \( (\kappa \sqrt{q - 1})^{|\gamma|} \), where \( \kappa \) is the probability that the spins on the SQ sublattice follow a fixed checkerboard pattern (say, 1 on the even sublattice and 2 on the odd sublattice) raised to the power 1/volume. This latter probability is easy to bound explicitly, yielding \( \kappa \leq [1 + v_{CG}(q,v)]^{-1}[q/(q - 5)]^{1/2} \), where \( v_{CG}(q,v) \) is the one for the corresponding unmodified lattice \( G_n \) or \( H_n \). This implies that, for all the lattices \( G'_n, \ldots, H'_n \), there is FLRO on the SQ sublattice whenever \( 0 \leq q \leq q_c(G_n) \sim O(1) \) [cf. Eq. (3)] and \( v \) is close to \(-1 \) (low temperature).

Let us also remark that the lattice \( G'_2 \) is covered by the general theory of Ref. \cite{18}, where it is proven that \( q_c > 3 \); moreover, a minor modification proves the same result for \( G'_n \) for all \( n \geq 2 \).

V. DATA FOR LATTICES \( G'_n \) AND \( H'_n \)

The lattices \( G'_n \) and \( H'_n \) for all \( n \) can be reduced to the union-jack (UJ) lattice (Fig. 3) with \( v_x = v \) and \( v_{\square} = v_{CG}(q,v) \) [cf. Eqs. (1) and (2) when \( v = -1 \); of course the same reduction holds for \( G_n \) and \( H_n \) by setting \( v_y = 0 \). We obtained high-precision estimates of the phase boundary of the UJ model in the \((v_x, v_{\square})\) plane by using the Jacobsen-Scullard (JS) method \cite{21} with untwisted square

TABLE I. Estimates of \( q_c(n) \) for the lattices \( G_n, G'_n, H_n, H'_n \) from the Jacobsen-Scullard (JS) method or the exact solution, and their large-\( n \) asymptote \( 2n/W(2n) \) from Eq. (3).

\begin{tabular}{|c|c|c|c|c|}
\hline
\( n \) & \( q_c(G_n) \) (exact) & \( q_c(G'_n) \) (JS) & \( q_c(H_n) \) (JS) & \( q_c(H'_n) \) (JS) & \( 2n/W(2n) \) (asympt.) \\
\hline
1 & 2.618034 & 3.74583(8) & 2.618034 & 3.74583(8) & 2.345751 \\
2 & 3.448678 & 4.48805(4) & 4 & 4.80794(5) & 3.327322 \\
4 & 4.942152 & 5.87902(5) & 5.617069 & 6.39269(4) & 4.981903 \\
8 & 7.565625 & 8.40372(3) & 8.30417 & 9.04238(2) & 7.792741 \\
16 & 12.164794 & 12.95103(1) & 12.93942 & 13.63211(2) & 12.621338 \\
64 & 34.667189 & 35.27672(3) & 35.482095 & 36.07477(3) & 35.780223 \\
\hline
\end{tabular}
bases of size up to $7 \times 7$ (294 edges) [30]. We checked these results for $q = 4, 8, 16, 32$ by Monte Carlo simulations using a cluster algorithm [31]. The estimated phase boundaries from both methods are shown in Fig. 4, along with the numerical estimates of $q_c$ from both methods. The estimates of $q_c(n)$ for the lattices $G_n, G_n', H_n, H_n'$ from the JS method (or the exact solution) are shown in Table I, where they are compared with the predicted large-$n$ asymptote $q_c(n) \approx 2n/W(2n)$ from Eq. (3). We then performed linear and quadratic extrapolations to locate the point $q_c$, where $v_c = -1$. The results are shown in Table II and Fig. 7 and agree well with the transfer-matrix estimates. For $q \gtrsim 8$ the specific heat diverges at the transition point like $L \sim q_c$, in agreement with the finite-size-scaling prediction for a first-order transition; for $4 < q < 8$ the transition is presumably also first order but with a large correlation length $\xi$, so that we are unable to observe the true $L \gg \xi$ asymptotic behavior.

VI. DATA FOR LATTICES $G_n''$ AND $H_n'''$

We studied the lattices $G_n''$ and $H_n'''$ for $n = 1, 2, 4, 8, 16, 32, 64$ (note that $G_1'' = \text{SQ}$ [6] and $H_1''' = \text{UJ}$ [20]) at $v = -1$, using transfer matrices with cylindrical boundary conditions on widths $L = 1, 2, 3, 4$ unit cells (Fig. 6). The computational complexity is linear in $n$. We estimated the location of the phase transition (which we expect to be first order whenever $q_c > 4$) using the crossings of the energies $E_{1/2}(v')$ [33]; the results are shown in Table II.

<table>
<thead>
<tr>
<th>$n$</th>
<th>$q_c(G_n'')$ (TM)</th>
<th>$q_c(H_n''')$ (TM)</th>
<th>$q_c(H_n''')$ (MC)</th>
<th>$2n/W(2n)$ (asympt.)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>3</td>
<td>4.31(3)</td>
<td>-</td>
<td>2.345751</td>
</tr>
<tr>
<td>2</td>
<td>3.63(2)</td>
<td>5.27(1)</td>
<td>5.26(2)</td>
<td>3.327322</td>
</tr>
<tr>
<td>4</td>
<td>5.02(1)</td>
<td>6.68(1)</td>
<td>6.67(3)</td>
<td>4.981903</td>
</tr>
<tr>
<td>8</td>
<td>7.60(1)</td>
<td>9.21(1)</td>
<td>9.21(7)</td>
<td>7.792741</td>
</tr>
<tr>
<td>16</td>
<td>12.18(2)</td>
<td>13.73(2)</td>
<td>13.73(10)</td>
<td>12.621338</td>
</tr>
<tr>
<td>32</td>
<td>20.29(3)</td>
<td>21.76(3)</td>
<td>21.76(32)</td>
<td>21.016077</td>
</tr>
<tr>
<td>64</td>
<td>34.70(5)</td>
<td>36.10(5)</td>
<td>36.14(8)</td>
<td>35.780223</td>
</tr>
</tbody>
</table>

For $H_n'''$ we checked these results by Monte Carlo. For three integer values of $q$ below the estimated $q_c$, we simulated the model at finite temperature and estimated the transition point $v_c(q)$. We then performed linear and quadratic extrapolations to locate the point $q_c$, where $v_c = -1$. The results are shown in Table II and Fig. 7 and agree well with the transfer-matrix estimates. For $q \gtrsim 8$ the specific heat diverges at the transition point like $L^{1/2}$, in agreement with the finite-size-scaling prediction for a first-order transition; for $4 < q < 8$ the transition is presumably also first order but with a large correlation length $\xi$, so that we are unable to observe the true $L \gg \xi$ asymptotic behavior.

VII. CONCLUSION

When a 2D model admits a height representation, it must be either critical or ordered. Until now criticality seemed to be the most common case, even though examples of order were known. But here we have exhibited several infinite families of 2D lattices (some of which are quadrangulations or Eulerian triangulations) in which the Potts antiferromagnet admitting a height representation ($q = 3$ or 4, respectively) is not only ordered but is in fact “arbitrarily strongly ordered” in the sense that $q_c$ is arbitrarily large. This unexpected result suggests that the prior belief may have things precisely backwards. Perhaps criticality is an exceptional case (arising, for instance, in situations with special symmetries) and order is to be generically expected. A key open question raised by this work is to understand why criticality arises when it does.

ACKNOWLEDGMENTS

This work was supported in part by NSFC Grants No. 10975127 and No. 11275185, the Chinese Academy of Sciences, French Grant No. ANR-10-BLAN-0414, the Institut Universitaire de France, Spanish MEC Grants No. FPA2009-08785 and No. MTM2011-24097, Czech GAČR Grant No. P201/12/2613, US NSF Grant No. PHY–0424082, and a computer donation from the Dell Corporation.
On the triangular lattice, both $q$-TWO-DIMENSIONAL POTTS ANTIFERROMAGNETS WITH A


Some exceptions are the constrained square-lattice four-state Potts antiferromagnet [9], the triangular-lattice antiferromagnetic spin-$s$ Ising model for large enough $s$ [C. Zeng and C. L. Henley, Phys. Rev. B 55, 14935 (1997)], the diced-lattice three-state Potts antiferromagnet [17], and the union-jack and bisected-hexagonal four-state Potts antiferromagnets [20], all of which appear to lie in a noncritical ordered phase at zero temperature.


On the triangular lattice, both $q = 2$ and $q = 4$ are critical at zero temperature and have height representations [14,34], but $q_c = 4$.


