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## LETTER TO THE EDITOR

**Collapsed two-dimensional polymers on a cylinder****Hsiao-Ping Hsu and Peter Grassberger**

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Online at [stacks.iop.org/JPhysA/35/L759](http://stacks.iop.org/JPhysA/35/L759)**Abstract**

Single partially confined collapsed polymers are studied in two dimensions. They are described by self-avoiding random walks with nearest-neighbour attractions below the  $\Theta$ -point, on the surface of an infinitely long cylinder. For the simulations we employ the pruned-enriched-Rosenbluth method (PERM). The same model had previously been studied for free polymers (infinite lattice, no boundaries) and for polymers on finite lattices with periodic boundary conditions. We verify the previous estimates of bulk densities, bulk free energies and surface tensions. We find that the free energy of a polymer with fixed length  $N$  has, for  $N \rightarrow \infty$ , a minimum at a finite cylinder radius  $R^*$  which diverges as  $T \rightarrow T_\theta$ . Furthermore, the surface tension vanishes roughly as  $(T_\theta - T)^\alpha$  for  $T \rightarrow T_\theta$  with  $\alpha \approx 1.7$ . The density in the interior of a globule scales as  $(T_\theta - T)^\beta$  with  $\beta \approx 0.32$ .

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Although the behaviour of polymers has been studied for many years, there are still a number of open questions. At high temperatures or in good solvents repulsive interactions (the excluded volume effect) dominate the conformation and the polymer chain tends to swell to a random coil. At low temperatures or in poor solvents attractive interactions between monomers dominate the conformation and the polymer chain tends to collapse and forms a compact dense globule. The coil–globule transition point is called the  $\Theta$ -point.

The behaviour in good solvents is accessible by field theory and thus well understood [1], e.g. the partition function scales as  $Z \sim \mu^N N^{\gamma-1}$  at  $T > T_\theta$ . In the collapsed phase, the fact that monomers are constrained to form a linear chain can be neglected in a first approximation, since they are anyhow kept close to each other by non-bonding attractive forces. Thus a collapsed polymer can essentially be viewed as a liquid droplet. Accordingly, Lifshitz mean field theory [2, 3] had predicted for  $T < T_\theta$  that there should be also a surface term,

$$Z(N, T) \sim a^N b^{-N^s} N^{\gamma-1} \quad (1)$$

with  $s = (d - 1)/d$  and  $b > 1$ . Note that no such surface term is expected in the random coil phase (i.e. for  $T > T_\theta$ ). First numerical indications for this were given by [4–6], and



Figure 1. Schematic drawing of a polymer chain growing on the surface of a cylinder.

theoretical arguments in favour of equation (1) were given in [7]. More recently, equation (1) was checked numerically in [8]. But it seems that there exists still no proof for equation (1), and it is not known how the surface tension (related to  $b$ ) scales for  $T \rightarrow T_\theta$ . Moreover, the exponent  $\gamma$  is not known [4, 6, 7], and it is not even clear whether it is universal in the collapsed phase.

In [8] we had studied two different geometries—globules on infinite lattices and polymers filling finite lattices with periodic boundary conditions—where surface tension should make very different contributions. In the present work we study a third geometry, namely polymers on the surface of an infinitely long cylinder, where the surface tension should make yet another contribution. Our main motivation was originally to have a simple geometry with straight ‘surfaces’ so that the surface tension and its  $T$ -dependence can be more easily measured. But we also found another interesting effect: the bulk fugacity  $\mu$  depends non-monotonically on the cylinder radius. It has a minimum at a finite radius  $R^*$  which diverges for  $T \rightarrow T_\theta$ . Thus a collapsed polymer, when wrapped around an elastic cylinder, will tend to squeeze it when  $R > R^*$  and widen it when  $R < R^*$ .

A collapsed polymer on the surface of an infinitely long cylinder is modelled as an interacting self-avoiding random walk (ISAW) on a square lattice with periodic boundary conditions in the vertical direction, as shown in figure 1. The vertical lattice size is  $h = 2\pi R$ . The collapse is mediated by nearest-neighbour attraction between monomers, described by a Boltzmann factor  $q = e^{-\beta\epsilon}$  for each pair of non-bonded monomers occupying nearest-neighbour sites on the lattice. Here  $\beta = 1/kT$  and  $\epsilon < 0$  is the attractive potential. The partition sum is therefore

$$Z = \sum_{\text{walks}} q^m \quad (2)$$

where  $m$  is the number of non-bonded occupied nearest-neighbour pairs. We are interested in polymers which form collapsed globules, therefore  $q > q_\theta$  with  $q_\theta = e^{-\epsilon/kT_\theta} \approx 1.95$  [10].

As argued in [2, 3] and verified in [4, 8], below the  $\Theta$ -point, the polymer forms a globule with uniform monomer density  $\rho$  for large  $N$  and  $\rho \rightarrow 0$  as  $T \rightarrow T_\theta$ . Thus, the free energy for a collapsed polymer with chain length  $N$  in infinite volume consists (except for a logarithmic term resulting from the last factor in equation (1)) of two parts: an extensive bulk contribution  $\propto N$  and a surface contribution  $\propto N^{(d-1)/d}$ . For  $d = 2$ , this reads (up to a factor  $1/T$  which will be suppressed in the following)

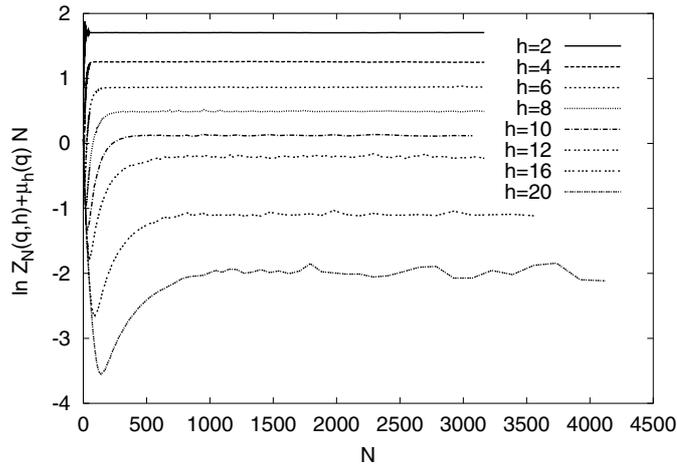
$$-\ln Z_N(q) \approx \mu_\infty(q)N + \tilde{\sigma}(q)\sqrt{N} + O(\ln N) \quad (3)$$

where  $\mu_\infty$  is the chemical potential per monomer in an infinite chain,  $\tilde{\sigma}$  is related to the surface tension (=free energy per unit of perimeter length)  $\sigma$  and to the monomer density  $\rho$  by

$$\sigma = \frac{1}{2}\tilde{\sigma}\sqrt{\frac{\rho}{\pi}}. \quad (4)$$

The latter follows by assuming the globule to be a disc with radius  $R_g$  fixed by  $N = \pi R_g^2 \rho$ .

Surface terms are absent from the free energy if one has a finite  $L \times L$  lattice with periodic boundary conditions, and if one chooses the polymer length such that the entire lattice is filled with finite density. More precisely, consider the quantity  $\ln Z_N(q) + \mu N$ . For  $T < T_\theta$  this is non-convex, and for  $\mu = \mu_L(q)$  it has two peaks of equal height, corresponding to the



**Figure 2.**  $\ln Z_N(q, h) + \mu_h(q)N$  versus  $N$  for  $q = 2.4$  and for various values of the cylinder circumference  $h$ . The values of  $\mu$  were determined by demanding that these curves become horizontal as  $N \rightarrow \infty$ . The peaks seen in the curves for  $h = 16$  and  $20$  are statistical fluctuations.

coexistence of two phases: one phase with very short chains and density  $\approx 0$  and another phase with density  $\rho_L(q)$ . For  $L \rightarrow \infty$ ,  $\mu_L(q)$  and  $\rho_L(q)$  tend monotonically [8] to  $\mu_\infty(q)$  and  $\rho$ .

Finally, on the surface of an infinitely long cylinder with perimeter  $h$ , a collapsed polymer forms a cylindrical blob with length  $L = \rho^{-1}N/h$  and surface free energy  $2\sigma h$ ,

$$-\ln Z_N(q, h) \approx \mu_h(q)N + 2\sigma(q)h. \quad (5)$$

Note that the chemical potential per monomer will in general depend on  $h$ , even in the limit  $N \rightarrow \infty$ . For  $h \rightarrow \infty$  we expect again to obtain the bulk limit,

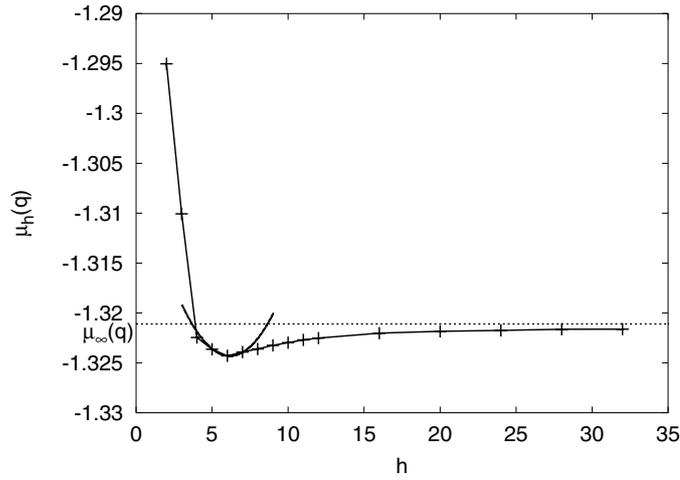
$$\mu_h(q) \rightarrow \mu_\infty(q) \quad h \rightarrow \infty. \quad (6)$$

In principle we should also expect  $\sigma(q)$  and  $\rho(q)$  to depend on  $h$ . But we shall not study the detailed behaviour of  $\rho$ , and we shall neglect the  $h$ -dependence of  $\sigma(q)$ , since the surface term is a small correction anyhow. We also might expect logarithmic terms in  $N$  to be present (just as also in the previous geometries), but they will also be neglected.

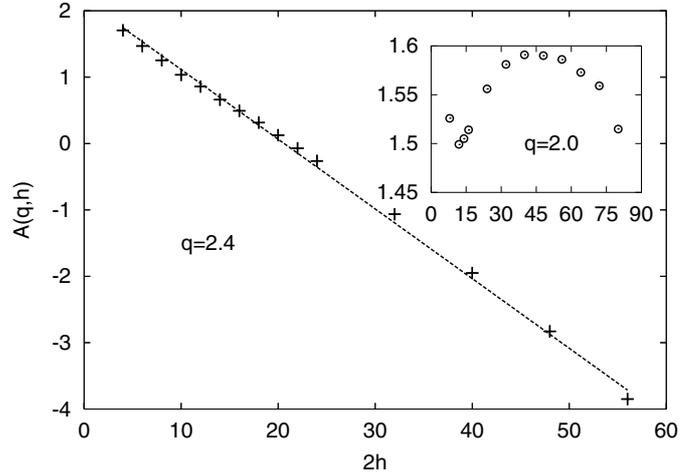
Our aim is to verify equation (5) numerically, to estimate the functions  $\mu_h(q)$  and  $\sigma(q)$  and to see whether they are compatible with the results for the other two geometries.

The pruned-enriched-Rosenbluth method (PERM) has been employed to study collapsed 2D polymers in [8, 10]. In [8] the predicted behaviour was verified numerically for globules on infinite lattices and for polymers filling finite lattices with periodic boundary conditions. More precisely, finite lattices were studied first, and from them estimates were obtained for  $\mu_\infty$  and  $\rho$  by extrapolation. Given these, subsequent simulations on infinite lattices gave rather precise estimates for  $\sigma(q)$ . But the latter was not done systematically for a wide range of temperatures, in particular the behaviour of  $\sigma(q)$  for  $q \rightarrow q_\theta$  was not studied. Thus we want to obtain also the latter in the present letter from cylinder geometry simulations.

For the first simulations, we choose  $q = 2.4$  which is deep inside the collapsed regime. From equation (5) we expect that plots of  $\ln Z_N(q, h) + \mu N$  against  $N$  should give horizontal lines for  $N \rightarrow \infty$ , if and only if  $\mu = \mu_h(q)$ . This gives estimates of  $\mu_h(q)$ . Using them, we plot in figure 2  $\ln Z_N(q, h) + \mu_h(q)N$  versus  $N$  for various values of the cylinder circumference  $h$ . For small  $h$  we find perfect agreement with our expectations. For larger  $h$  we see that PERM has problems in sampling long chains correctly, as should be expected. Fluctuations similar



**Figure 3.**  $\mu_h(q = 2.4)$  as obtained from figure 2, plotted against  $h$ . There is a minimum at  $h = 6.13 \pm 0.01$  which is determined by fitting a parabola (dashed curve) going through the three lowest points (which are at  $h = 5, 6$  and  $7$ ). The dotted line is the asymptote  $\mu_\infty(q) = -1.3211(1)$  approached for  $h \rightarrow \infty$ .



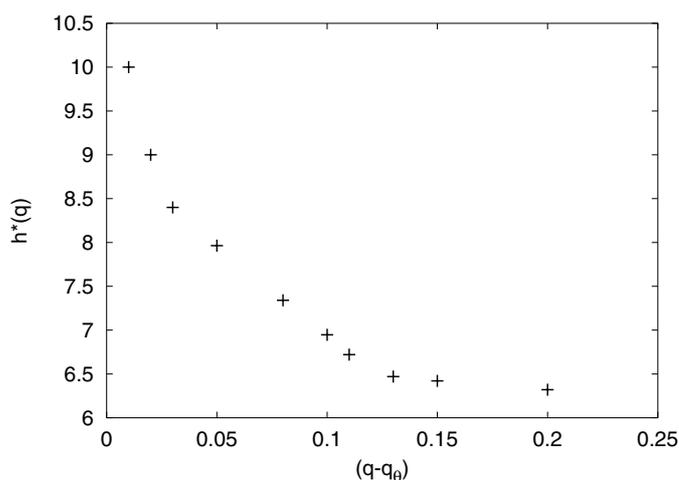
**Figure 4.**  $A(q, h)$  versus  $2h$  for  $q = 2.0$  and  $2.4$ . The dashed line has slope  $-\sigma = -0.105$ . The points for  $h = 16$  and  $h = 20$  have large errors for  $q = 2.4$ , as also seen from figure 2. Results for  $q = 2.0$  (inset) show the difficulty of measuring surface tension near  $\Theta$ -point.

to those seen in the curve for  $h = 16$  prevent us from going to much larger  $h$ , except for  $q$  close to  $q_\theta$  where the efficiency of PERM increases.

Values of  $\mu_h(q)$ , again for  $q = 2.4$ , are plotted in figure 3. For large  $h$ , we obtain  $\mu_\infty(q) = -1.3211(1)$  in perfect agreement with [8]. But, surprisingly, the dependence of  $\mu_h(q)$  on  $h$  is not monotonic. Instead, we see a minimum at a finite value of  $h$ . In our simulations  $h$  is an integer, but we can interpolate to continuous values by fitting a parabola,  $\mu_h(q) = \mu_{\min} + b(h - h^*)^2$ , through the three lowest points, as shown in figure 3. Such a fit gives  $h^*(q = 2.4) = 6.13(1)$  and  $\mu_{\min}(q = 2.4) = -1.3243(1)$ .

Finally, in order to obtain the surface tension, we plot in figure 4

$$A(q, h) = \lim_{N \rightarrow \infty} [\ln Z_N(q, h) + \mu_h(q)N] \quad (7)$$



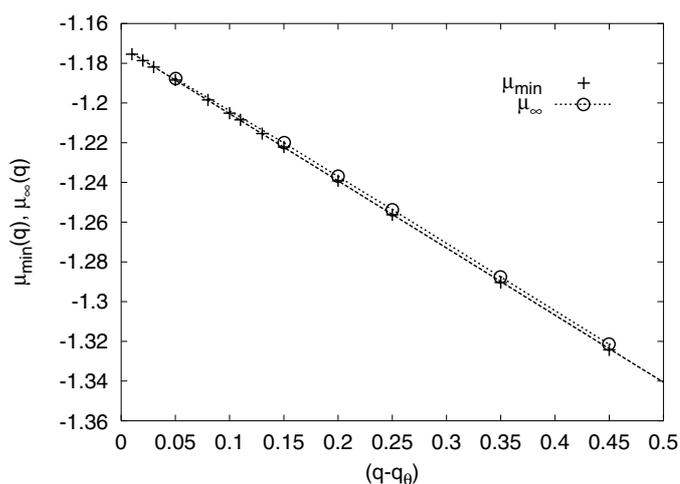
**Figure 5.** ‘Optimal’ cylinder circumference  $h^*(q)$  against  $(q - q_\theta)$ . At  $h = h^*(q)$ , the chemical potential is minimal,  $\mu_{h^*}(q) = \mu_{\min}(q)$ .

against  $2h$ . The limit is of course obtained by using the parts of the curves in figure 2 where they are horizontal. From equation (5) we expect this to give a straight line with slope  $-\sigma$ . The result shows that  $\sigma = 0.105(5)$  at  $q = 2.40$ . Actually, the line is slightly curved, showing that systematic errors are not negligible. For large  $h$  they seem to arise from insufficient sampling. For small  $h$  they result from the neglected terms discussed above. The present estimate should be compared to the estimate  $\sigma = 0.119(5)$  from [8]. Although the discrepancy is about 2 standard deviations, we consider this as reasonable agreement. Results for  $q = 2.0$  are also shown in figure 4 and illustrate the large corrections near the  $\Theta$ -point.

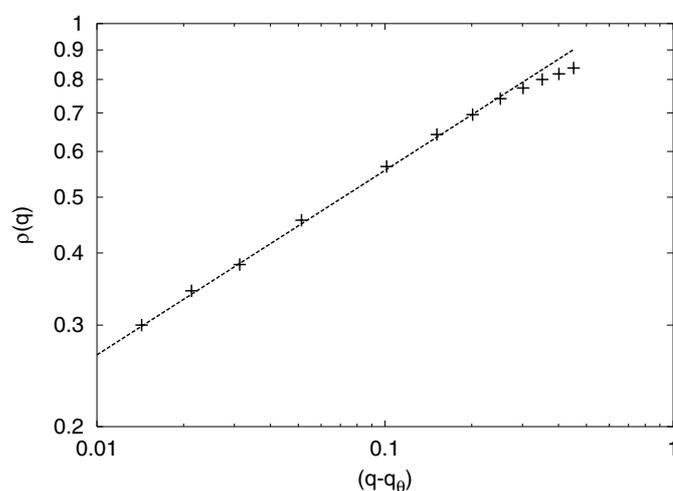
Analogous simulations as shown in figures 2–4 were also done at other values of  $q$ . The lengths  $N$  depended strongly on  $q$ . Close to the  $\Theta$ -point the algorithm is very efficient and we obtained reliable high statistics data for  $N$  up to 15 000. The more the polymer is collapsed, the more difficult are the simulations. We stopped at  $q = 2.4$ . As seen from figure 2, the errors decrease when  $h$  becomes small, most likely because the walk is effectively Markovian with memory  $\sim h^2$ , and PERM is most efficient when memory is short. In all those simulations the results were similar.  $\mu_h(q)$  had a minimum at a finite  $h^*(q)$ , and  $A(q, h)$  versus  $h$  was roughly linear with a negative slope. Results of  $h^*(q)$ ,  $\mu_{\min}(q)$ ,  $\rho(q)$  and  $\sigma(q)$  for a wide range of  $q$  are shown in figures 5–8.

In figure 5 we can see that  $h^*(q)$  increases as  $q$  decreases. This is to be expected since the  $\Theta$  transition is continuous, with a diverging correlation length for  $T \rightarrow T_\theta$ . Apart from the lattice constant (which is set to 1 here), the correlation length is the only relevant length scale. The divergence of  $h^*(q)$  when approaching the  $\Theta$ -point thus indicates that the existence of the minimum is not a lattice artefact and is presumably a genuine feature of the continuum theory. The simulations are not precise enough to allow a meaningful fit to a power law and a determination of a critical exponent. This exponent would also be affected strongly by the uncertainty of  $T_\theta$ . To reduce the latter to a minimum, we performed extensive simulations of free (i.e. infinite lattice) ISAWs near the estimated  $\Theta$ -point, with the result  $q_\theta = 1.9487(5)$  which is in good agreement with the best previous estimate [10].

In figure 6 we plot  $\mu_{\min}(q)$  and  $\mu_\infty(q)$  versus  $(q - q_\theta)$ . We see that both decrease linearly with temperature. The difference between them depends weakly on the temperature. It seems to vanish for  $T \rightarrow T_\theta$ .



**Figure 6.**  $\mu_{\min}$  and  $\mu_{\infty}$  plotted against  $(q - q_{\theta})$  for various values of  $q$ . The slope of the straight line is roughly equal  $-0.34$ .  $\mu_{\infty}$  were obtained from extrapolation to large  $h$ .



**Figure 7.** log–log plot of  $\rho$ , as obtained from the growth of  $\langle |x_N| \rangle$  with  $N$  for large  $N$  and  $h$ , against  $(q - q_{\theta})$ . For this plot we used  $q_{\theta} = 1.9487$ . The dashed line is  $\rho \propto (q - q_{\theta})^{0.32}$ .

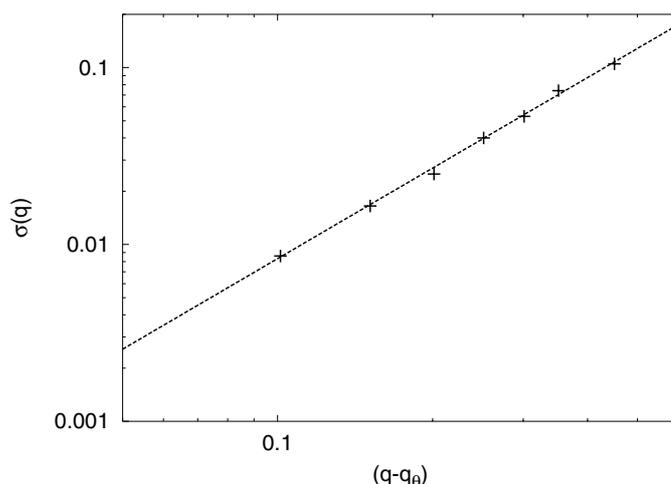
The monomer density  $\rho$  inside the ‘globule’ is constant for  $N \gg h^2 \gg 1$  and fixed temperature. It can be measured by the  $x$ -coordinate (parallel to the cylinder) of the end point of the chain. Since the chain has to grow either in the positive or in the negative direction, we have

$$\frac{d\langle |x_N| \rangle}{dN} \sim (h\rho)^{-1} \quad N \rightarrow \infty. \quad (8)$$

For small values of  $h$  there are finite-size corrections to the density, but it is easy to extrapolate to  $h \rightarrow \infty$ . The resulting densities are plotted versus  $q$  in figure 7. They show a power law

$$\rho \sim (q - q_{\theta})^{\beta} \quad (9)$$

with  $\beta = 0.32 \pm 0.02$ . We are not aware of previous estimates of  $\beta$ , either theoretically or by simulations.



**Figure 8.** log–log plot of  $\sigma$ , as obtained from figure 4, against  $(q - q_\theta)$ . For this plot we used  $q_\theta = 1.9487$ . The dashed line is  $\sigma \propto (q - q_\theta)^{1.7}$ .

The values of the surface tension  $\sigma(q)$  shown in figure 8 are obtained in the same way as shown in figure 4. They follow a power law,  $\sigma \sim (q - q_\theta)^\alpha$ , with  $\alpha \approx 1.7 \pm 0.1$ . For the parameter  $\tilde{\sigma}$  this implies that

$$\tilde{\sigma} \sim (q - q_\theta)^{\alpha - \beta/2} \sim (q - q_\theta)^{1.54 \pm 0.1}. \quad (10)$$

This is incompatible with the theoretical prediction  $\tilde{\sigma} \sim (T_\theta - T)^{7/6}$  quoted in [6].

In summary, we have studied collapsed polymers on the surface of a cylinder. Using the PERM algorithm we could simulate chains of lengths between 3000 (far below  $T_\theta$ ) and  $10^4$  (near  $T_\theta$ ) for a wide range of temperatures. We have verified that the concept of a surface tension applies to 2D collapsed polymers and have determined its scaling near the  $\Theta$ -point. We also determined the scaling of the monomer density in the collapsed phase, as  $T \rightarrow T_\theta$ . In addition, we found a surprising non-monotonic dependence of the chemical potential on the perimeter of the cylinder. It shows that for any temperature there is a special value of the perimeter where the free energy of the polymer is minimal. The fact that this value diverges when the  $\Theta$ -temperature is approached suggests that this is not an artefact of the square lattice but a generic feature.

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### References

- [1] de Gennes P G 1979 *Scaling Concepts in Polymer Physics* (Ithaca, NY: Cornell University Press)
- [2] Grosberg A Yu and Kuznetsov D V 1992 *Macromolecules* **25** 1970
- [3] Lifshitz I M 1968 *JETP* **28** 1280
- [4] Owczarek A L, Prellberg T and Brak R 1993 *Phys. Rev. Lett.* **70** 951

- [5] Grassberger P and Hegger R 1995 *J. Chem. Phys.* **102** 6881
- [6] Nidras P P 1996 *J. Phys. A: Math. Gen.* **29** 7929
- [7] Duplantier B 1993 *Phys. Rev. Lett.* **71** 4274
- [8] Grassberger P and Hsu H-P 2002 *Phys. Rev. E* **65** 031807
- [9] Grassberger P 1997 *Phys. Rev. E* **56** 3682
- [10] Barkema G T, Bastolla U and Grassberger P 1998 *J. Stat. Phys.* **90** 1311