Critical adsorption of polymers in a medium with long-range correlated quenched disorder.

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We investigated the influence of short- and long-range correlated quenched disorder introduced into the medium on the process of adsorption of long-flexible polymer chains on a wall by using the field theoretical approach in $d = 4 - \epsilon$ and directly in $d = 3$ dimensions up to one-loop order for the semi-infinite $|\phi|^4$ m-vector model (in the limit $m \to 0$) with a boundary. This allows us to describe the critical behavior of long-flexible polymer chains in the vicinity of the surface and to obtain the whole set of surface critical exponents at the special surface transition ($c = c_{ads}$), which separates the nonadsorbed region $c > c_{ads}$ from the adsorbed one $c < c_{ads}$. In the case of very strong correlation of the disorder we obtained that a polymer collapses and then adsorbs on a wall. The obtained results indicate that for the systems with long-range-correlated quenched disorder the new set of surface critical exponents arises and show that polymer chains in solution with that type of disorder adsorb on a wall stronger than polymer chains in pure medium.

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The critical behavior of the real systems with different kinds of disorder is of considerable theoretical and experimental interest. In common, the real physical systems are usually characterized by the presence of different kinds of defect and impurities that may be localized inside the bulk or at the boundary. The present work is connected with investigation of adsorption phenomena of long-flexible polymer chains in disordered media with correlated structural defects. The main question is: does a small amount of long-range correlated quenched disorder introduced into the medium induce changes to the universal properties of long-flexible polymer chains at the adsorption on a wall?

As it is known, long flexible polymer chains in a good solvent are perfectly described by a model of self-avoiding walks (SAW) on a regular lattice. Their scaling properties in the limit of infinite number of steps may be derived by a formal $m \to 0$ limit of the $O(m)$ vector model at its critical point. Thus for the average square end-to-end distance $R$ and the number of configurations $Z_N$ of a SAW with $N$ steps and one end fixed or with both ends fixed on a regular lattice in the asymptotic limit $N \to \infty$ the following relations take place

$$< R^2 > \sim N^{2\nu}, \quad Z_N \sim q^N N^{\gamma-1}, \quad Z_N(r) \sim q^N N^{-(2-\alpha)},$$

where $\nu$ and $\gamma$ are the universal correlation length and susceptibility exponents for the $m = 0$ model, $d$ is the space dimensionality, $q$ is a nonuniversal fugacity and $r^2 = (\vec{r}_A - \vec{r}_B)^2$. What happens when disorder will be introduced into the system? As it is known the infinite systems are described by the Landau-Ginzburg-Wilson Hamiltonian

$$H = \int d^d x \left[ \frac{1}{2} \left| \nabla \phi \right|^2 + \frac{1}{2} (m_i^2 + \delta \tau(x)) \left| \phi_i \right|^2 + \frac{1}{4!} v_0 (\phi^2)^2 \right],$$

where $\phi(x)$ is an $m$-vector field with the components $\phi_i(x)$, $i = 1, ..., m$. The inhomogeneities in the system cause fluctuations in the local transition temperature. Here $m_i^2$ is the "bare mass", which can be interpreted in this case as chemical potentials for monomers in the bulk, $\delta \tau(x)$ represents the quenched random-temperature disorder, with $< \delta \tau(x) > = 0$ and

$$\frac{1}{8} < \delta \tau(x) \delta \tau(x') > = g(|x|),$$

where angular brackets $< ... >$ denote configurational averaging over quenched disorder. The pair correlation function $g(|x|)$ is assumed to fall off with the distance as

$$g(|x|) \sim \frac{1}{x^a}$$

for large $x = (r, z)$, where $a$ is a constant and $x = |x|$.

The Fourier-transform $\tilde{g}(k)$ of $g(x)$ for small $k$ is
\[ \hat{g}(k) \sim u_0 + w_0 |k|^{a-d}. \]  

This corresponds to the so-called long-range-correlated "random-temperature" disorder. In the case of random uncorrelated pointlike (or short-range-correlated) disorder the site-occupation correlation function is \( g(x) \sim \delta(x) \) and its Fourier-transform

\[ \hat{g}(k) \sim u_0. \]  

Applying the replica method in order to average the free energy over different configurations of quenched disorder, it is possible to construct an effective Hamiltonian of the \(|\phi^4| m\)-vector model with a long-range-correlated disorder

\[
H_{eff} = \sum_{\alpha=1}^{n} \int_{V} d^{d}x \left[ \frac{1}{2} |\nabla \phi_{\alpha}|^2 + \frac{1}{2} m_{0}^{2} \phi_{\alpha}^2 + \frac{1}{4!} v_0 (\phi_{\alpha}^2)^2 \right. \\
- \left. \sum_{\alpha,\beta=1}^{n} \int_{V} d^{d}x_1 d^{d}x_1' g(|x_1-x_1'|) \phi_{\alpha}^2(x_1) \phi_{\beta}^2(x_1'). \right]
\]  

We are interested in investigation of the polymer limit \( m \to 0 \) of such kind of model which can be interpreted as a model of long-flexible polymer chains in a disordered medium. Here Greek indexes denote replicas, and the replica limit \( n \to 0 \) is implied. If \( a \geq d \), then \( w_0 \) term irrelevant. This corresponds to random uncorrelated pointlike disorder (or short-range-correlated random disorder). As noticed by Kim\(^3\), in this case in the limit \( m, n \to 0 \) both \( v_0 \) and \( w_0 \) terms are of the same symmetry. It indicates that a weak quenched uncorrelated disorder is irrelevant for SAWs\(^4\).

Thus, it is impossible to "naively" apply the Harris criterion\(^2\) to the SAW problem, though the critical exponent \( \alpha \) of a SAW on the \( d = 3 \) dimensional pure lattice is positive (\( \alpha(d = 3) = 0.235 \pm 0.003\)\(^8\)).

If \( a < d \), the term \( w_0 k^{a-d} \) is relevant and the long-range-correlated disorder is relevant for SAWs (see\(^7\)).

The presence of a surface or a wall leads to the appearance of additional problems in analysis of the critical behavior. The investigation of adsorption phenomena of long-flexible polymers on a hard wall was a subject of a series of works (for the sake of brevity we notice only few of them\(^2,8,9\)). It should be mentioned that in the case of systems with a surface, the different kinds of defects and impurities may be localized inside the bulk or at the boundary.

As was found in\(^10\), introducing into the system short-range correlated random quenched surface disorder is irrelevant for critical behavior, but long-range correlated quenched surface disorder with \( g(r) \sim \frac{1}{r} \) and with the corresponding Fourier transform

\[ \hat{g}(p) \sim u_0 + w_0 p^{a-d+1} \]  

can be relevant, if \( a < d - 1 \) and is irrelevant if \( a \geq d - 1 \).

Our present work is connected with the investigation of the adsorption phenomena on a planar wall of long flexible polymer chains inserted into disordered medium with long-range-correlated quenched disorder. The solution of polymer chains in this semi-infinite space is sufficiently dilute, so that interchain interactions can be neglected and it is sufficient to consider surface effects for configurations of a single chain.

We pay particular attention to the investigation of special surface transition \( c = c_{ads} \), which separates the nonadsorbed region \( c > c_{ads} \) from the adsorbed one \( c < c_{ads} \). In the latter two regions, the chemical potential \( c \) for monomers on the surface does not stay fixed but tends to \(+\infty \) and \(-\infty \), respectively\(^8,9\).

The effective Hamiltonian of the semi-infinite \(|\phi^4| m\)-vector model with long-range-correlated disorder in this case is

\[
H_{eff} = \sum_{\alpha=1}^{n} \int_{V} d^{d}x \left[ \frac{1}{2} |\nabla \phi_{\alpha}|^2 + \frac{1}{2} m_{0}^{2} \phi_{\alpha}^2 + \frac{1}{4!} v_0 (\phi_{\alpha}^2)^2 \right. \\
- \left. \sum_{\alpha,\beta=1}^{n} \int_{V} d^{d}x_1 d^{d}x_1' g(|x_1-x_1'|) \phi_{\alpha}^2(x_1) \phi_{\beta}^2(x_1') + \frac{c_0}{2} \sum_{\alpha=1}^{n} \int_{V} d^{d-1}r \phi_{\alpha}^2(r, z = 0). \right]
\]  

In accordance with the fact that we have to deal with semi-infinite geometry \( x = (r, z \geq 0) \), only parallel Fourier transformations in \( d - 1 \) dimensions will be performed. The parallel Fourier transform \( \hat{g}(q, z) \) of (4) is

\[ \hat{g}(q, z) \sim u_0 \frac{2^{2d+1}}{\Gamma(\frac{d-a}{2})} \frac{\pi^{\frac{d-a-1}{2}}}{z^{\frac{d-a-1}{2}}} K_{\frac{d-a+1}{2}}(qz), \]
where $K_{d-1}(qz)$ is the modified Bessel function and $q = |q|$, where $q$ is $d-1$ dimensional vector. In the case of small $q$ and $z$ we obtain the relation

$$
\tilde{g}(q, z) \sim u_0 + w_0 q^{a-d+1} + w_0' z^{d-a-1},
$$

(11)

which agrees with the predictions obtained in\textsuperscript{10} (see (8) in the case $z = 0$). In the general case of arbitrary $z$ (from $z = 0$ on the wall to $z \to \infty$) we must take into account the Fourier transform $\tilde{g}(q, z)$ of the form (10).

The corresponding correlation function of the model (9), which involves $N$ fields $\phi(x_i)$ at distinct points $x_i (1 \leq i \leq N)$ in the bulk and $M$ fields $\phi(r_j, z = 0) \equiv \phi_s(r_j)$ at distinct points on the wall with parallel coordinates $r_j (1 \leq j \leq M)$, has the form

$$
G^{(N,M)}(\{x_i\}, \{r_j\}) = \langle \prod_{i=1}^{N} \phi(x_i) \prod_{j=1}^{M} \phi_s(r_j) \rangle .
$$

(12)

The corresponding full free propagator in the mixed $pz$ representation is given by

$$
G(p, z, z') = \frac{1}{2\kappa_0} \left[ e^{-\kappa_0|z-z'|} - \frac{c_0 - \kappa_0}{c_0 + \kappa_0} e^{-\kappa_0(z+z')} \right],
$$

(13)

where $\kappa_0 = \sqrt{p^2 + m_0^2}$ with $p$ being the value of the parallel momentum $p$ associated with the $d-1$ translationally invariant directions in the system.

There are two possible cases: a) when two ends of polymer are attached to the wall (in such a case we have to deal with a calculation of two point correlation function $G^{(0,2)}(r, z = 0; r', z' = 0)$), and b) when one end of the polymer is in the bulk and the other one is attached to the wall (specific $G^{(1,1)}(x; r', z' = 0)$). In order to obtain the universal surface critical exponents characterizing the adsorption of long-flexible polymer chains inserted into the medium with long-range-correlated quenched disorder on the wall, it is sufficient to consider the correlation function of two surface fields $G^{(0,2)}(r, z = 0; r', z' = 0)$. The universal surface critical exponents for such systems depend on the dimensionality of space $d$, the value of order parameter components $m(m \to 0)$ and the range of the correlations, i.e. the value of $a$.

As it is known, in the theory of semi-infinite systems the bulk field $\phi(x)$ and the surface field $\phi_s(r)$ should be reparameterized by different uv-finite renormalization factors $Z_\phi(u, v, w)$ and $Z_1(u, v, w)$,

$$
\phi(x) = Z_\phi^{1/2} \phi_R(x) \quad \text{and} \quad \phi_s(r) = Z_1^{1/2} Z_1^{1/2} \phi_{s,R}(r).
$$

The renormalized correlation function involving $N$ bulk and $M$ surface fields can be written as

$$
G_R^{(N,M)}(p; m, u, v, w, c) = Z_\phi^{-(N+M)/2} Z_1^{-M/2} G^{(N,M)}(p; m_0, u_0, v_0, w_0, c).
$$

(14)

It should be mentioned that the typical bulk short-distance singularities, which are present in the correlation function $G^{(0,2)}$, can be subtracted after performing the mass renormalization. For distinguished the parallel and perpendicular directions we obtain:

$$
m_0^2 = m^2 - t_1^{(0)} I_1(m^2) + t_2^{(0)} I_2(m^2),
$$

(15)

where

$$
t_1^{(0)} = \frac{1}{3} (v_0 - u_0 - \frac{w_0 m^{a-d}}{\cos[\pi (a-d)/2]}), \quad t_2^{(0)} = \frac{w_0}{3 \sqrt{\pi}} \frac{\Gamma(d-a-1)}{\Gamma(d-a/2)},
$$

(16)

and

$$
I_1(m^2) = \frac{1}{(2\pi)^{d-1}} \int \frac{d^{d-1} q}{2\kappa_q} q^{a-d+1} 2 F_1 \left[ \frac{1}{2}, 1, \frac{3+a-d}{2}, \frac{q^2}{2\kappa_q^2} \right]
$$

(17)

with $\kappa_q = \sqrt{q^2 + m^2}$ and

$$
I_2 = \frac{1}{(2\pi)^{d-1}} \int d^{d-1} q \frac{\sqrt{|q| a-d+1} 2 F_1 \left[ \frac{1}{2}, 1, \frac{3+a-d}{2}, \frac{q^2}{2\kappa_q^2} \right]}{2\kappa_q^2},
$$

where $2 F_1$ is the hypergeometric function.
According to the above mentioned notations, we have in the effective Hamiltonian only two coupling constants, \( V_0 = v_0 - u_0 \) and \( w_0 \) (we keep notation \( v_0 \) for \( V_0 \)).

The renormalized coupling constants \( v, w \) are fixed via the standard normalization conditions of the infinite-volume theory\(^7\)

\[
m^4 - d v = \Gamma^{(4)}_{b, R, v}(\{ q \}; m^2, v, w)|_{q=0},
\]

\[
m^4 - a w = \Gamma^{(4)}_{b, R, w}(\{ q \}; m^2, v, w)|_{q=0},
\]

(18)

where \( \Gamma^{(4)}_{b, R, v} \) and \( \Gamma^{(4)}_{b, R, w} \) are the \( v \) and \( w \) term symmetry contributions to the four-point vertex function. To the present accuracy of calculation at one-loop order, the vertex renormalization gives: \( v = v_0 m^{d-4} \) and \( w = w_0 m^{a-4} \).

In order to remove the short-distance singularities of the correlation function \( G^{(0,2)} \) located in the vicinity of the surface, the surface-enhancement shift \( \delta c \) is required. In accordance with this, the new normalization condition should be introduced for the definition of the surface-enhancement shift and surface renormalization factor \( Z_1 \). By analogy with magnetic systems\(^{11,12,13}\), the renormalized surface two-point correlation function in our case is normalized in such a manner\(^11\) that at zero external momentum it should coincide with the lowest order perturbation expansion of the surface susceptibility \( \chi_s(p) = G^{(0,2)}(p) \)

\[
G^{(0,2)}(p; m_0, v_0, w_0, c_0) = \frac{1}{c_0 + \sqrt{p^2 + m_0^2}} + O(v_0, w_0)
\]

(19)

and its first derivatives with respect to \( p^2 \). Thus we obtain the necessary surface normalization conditions

\[
G^{(0,2)}_R(0; m, v, w, c) = \frac{1}{m + c}
\]

(20)

and

\[
\frac{\partial G^{(0,2)}_R(p; m, v, w, c)}{\partial p^2}
\]

\[
|_{p=0} = \frac{1}{2m(m + c)^2}.
\]

(21)

Equation (20) defines the required surface-enhancement shift \( \delta c \) and shows that the surface susceptibility diverges at \( m = c = 0 \). This point corresponds to the multicritical point \( (m^2_{sc}, c^*_{sc}) \) at which the special transition takes place.

From the normalization condition of Eq. (21) and expression for the renormalized correlation function of Eq. (14), we can define the renormalization factor \( Z_\parallel = Z_1 Z_\phi \) by

\[
Z_\parallel^{-1} = 2m \frac{\partial}{\partial p^2} [G^{(0,2)}(p)]^{-1} |_{p^2=0} = \lim_{p^2 \rightarrow 0} m \frac{\partial}{\partial p} [G^{(0,2)}(p)]^{-1}.
\]

(22)

Asymptotically close to the critical point the two point renormalized correlation functions \( G^{(N,M)}_R \) satisfy the corresponding homogeneous Callan-Symanzik equations\(^{11,12}\)

\[
\left[ m \frac{\partial}{\partial m} + \beta_v(v, w) \frac{\partial}{\partial v} + \beta_w(v, w) \frac{\partial}{\partial w} + \frac{N + M}{2} \eta(v, w) \right. \\
+ \frac{M}{2} \eta_\phi^{sp}(v, w) \right] G^{(N,M)}_R(0; m, v, w, c) = 0,
\]

(23)

where the \( \beta \)-functions are \( \beta_v(v, w) = m \frac{\partial}{\partial m} \big|_{LR} v \), \( \beta_w(v, w) = m \frac{\partial}{\partial m} \big|_{LR} w \), the exponents \( \eta \) and \( \eta_\phi^{sp} \) are

\[
\eta = \frac{m}{\partial m} \ln Z_\phi \big|_{LR}, \quad \eta_\phi^{sp} = \eta \frac{\partial}{\partial m} \ln Z_1 \big|_{LR},
\]

(24)

and where LR is the long-range fixed point. The simple scaling dimensional analysis of \( G^{(0,2)}_R \) and mass dependence of \( Z \) factors, defines the surface correlation exponent \( \eta_\parallel^{sp} \) via

\[
\eta_\parallel^{sp} = \eta_\phi^{sp} + \eta.
\]

(25)
From Eqs. (22),(24) and (25), we obtain for the surface correlation exponent \( \eta_{sp}^{\parallel} \)

\[
\eta_{sp}^{\parallel} = n \frac{\partial}{\partial m} \ln Z_{\parallel} \bigg|_{LR} = \beta_v(v, w) \frac{\partial \ln Z_{\parallel}(v, w)}{\partial v} + \beta_w(v, w) \frac{\partial \ln Z_{\parallel}(v, w)}{\partial w} \bigg|_{LR},
\]

where the one-loop pieces of the \( \beta \) functions are

\[
\beta_v(\bar{v}, \bar{w}) = -\bar{v} + \bar{v}^2 - (3 f_1(a) - f_2(a)) \bar{w},
\]

\[
\beta_w(\bar{v}, \bar{w}) = -(4 - a) \bar{w} - (f_1(a) - f_2(a)) \bar{w}^2 + \frac{\bar{v} \bar{w}}{2}.
\]

The renormalized coupling constants \( v \) and \( w \) are normalized in a standard fashion so that

\[
\bar{v} = \frac{4}{3} v I_1, \quad \bar{w} = \frac{4}{3} w I_1,
\]

and the integral \( I_1 \) (see [17]) in the case of \( d = 3 \) is equal to \( 1/8 \pi \) and in the case of \( d = 4 - \epsilon \) it is \( I_1 = 2^{-d/2} \Gamma(\epsilon/2) \)

( the case \( \epsilon = 4 - d \) and \( \delta = 4 - a \) is presented in Appendix 1). The coefficients \( f_i(a) \) expressed via the one-loop integrals represented in \( 7 \) and are \( 14 \)

\[
\begin{align*}
  f_1(a) &= \frac{(a - 2)(a - 4)}{2 \sin(\pi a/2)}, \\
  f_2(a) &= \frac{(a - 2)(a - 3)(a - 4)}{48 \pi \sin(\pi(a/2 - 1))}, \\
  f_3(a) &= \frac{(2a - 5)(2a - 7)}{2 \sin(\pi(a - 3/2))}.
\end{align*}
\]

After performing the integration of the corresponding Feynman integrals in the renormalized two point correlation function \( G^{(0, 2)} \) we obtain for the renormalization factor \( Z_{\parallel} \) at one-loop order (in the case \( d = 3 \) ) the following result

\[
Z_{\parallel} = 1 + \frac{\bar{v}}{8} - \frac{\bar{w}}{8} \frac{2^{(a-3)/2}}{\cos(\pi(a - 3)/2)} (2^{(a-1)/2} - 2^{(a+1)/2} - \frac{2^{(a-1)/2}}{a - 5}).
\]

Combining the renormalization factor \( Z_{\parallel} \) together with one-loop pieces of the \( \beta \)-functions, according to Eq.(26) we finally obtain the following expression for the surface critical exponent \( \eta_{sp}^{\parallel} \)

\[
\eta_{sp}^{\parallel} = -\frac{\bar{v}}{8} + \frac{\bar{w}}{8} \frac{2^{(a-3)/2}(4 - a)}{\cos(\pi(3 - a)/2)} (2^{(a-1)/2} - 2^{(a+1)/2} - \frac{2^{(a-1)/2}}{a - 5}).
\]

The series for other surface critical exponents can be calculated on the basis of surface scaling relations (see Appendix 2) and one-loop series for bulk critical exponents

\[
\begin{align*}
  \nu^{-1} &= 2 - \frac{\bar{v}}{4} + \frac{f_1(a) - f_2(a)}{2} \bar{w}, \\
  \eta &= \frac{1}{2} f_2(a).
\end{align*}
\]

The results of our calculation are presented in Table 1.

As it is easy to see, in the case \( a = d = 3 \), which corresponds to random uncorrelated pointlike disorder (or short-range-correlated disorder) the obtained one-loop results for the surface critical exponents coincide with results for pure model (see\(^1\)\(^1\), Padé aproximants \([1/0]\)). Besides, for the case of a medium with long-range-correlated quenched disorder, the process of adsorption of long-flexible polymer chains on a wall is characterized by the new class of surface
Thus, for the renormalization factor $\epsilon$ in obtain for the surface critical exponent $z$ (distance and the value of the mean number of polymer chain free ends in a layer between distances $a$ order calculated at the pure (the case $a = 3$ with $(v^* = 1.632, w^* = 0)$) and LR stable fixed point for different fixed values of the correlation parameter, $2 < a < 3$.

| $a$  | $v^*$ | $w^*$ | $\eta_\parallel$ | $\eta_{\perp}$ | $\beta_1$ | $\gamma_1$ | $\gamma_{11}$ | $\delta_1$ | $\delta_{11}$ |
|------|-------|-------|-------------------|----------------|----------|------------|-------------|----------|-------------|
| 3.0  | 1.63  | 0.00  | -0.204           | -0.102         | 0.250    | 0.704      | 1.255       | 6.020    | 3.816       |
| 2.9  | 4.13  | 1.47  | -0.716           | -0.358         | 0.154    | 1.024      | 1.511       | 8.581    | 5.865       |
| 2.8  | 4.73  | 1.68  | -0.843           | -0.421         | 0.134    | 1.111      | 1.590       | 9.213    | 6.372       |
| 2.7  | 5.31  | 1.81  | -0.967           | -0.483         | 0.117    | 1.200      | 1.675       | 9.832    | 6.867       |
| 2.6  | 5.89  | 1.87  | -1.094           | -0.546         | 0.100    | 1.294      | 1.768       | 10.467   | 7.376       |
| 2.5  | 6.48  | 1.89  | -1.235           | -0.616         | 0.081    | 1.398      | 1.869       | 11.173   | 7.941       |
| 2.4  | 7.10  | 1.87  | -1.404           | -0.700         | 0.058    | 1.520      | 1.985       | 12.017   | 8.617       |
| 2.3  | 7.76  | 1.84  | -1.638           | -0.816         | 0.019    | 1.676      | 2.122       | 13.183   | 9.551       |

Critical exponents. They allow us to find useful characteristics which describe the critical behaviour of long-flexible polymer chains at the adsorption on a wall.

Since the bulk translational invariance is broken near the surface, the number of configurations $Z_{\parallel}$ depends on the distance $z_A = z$ of the starting point of the chain from the surface and behaves as $Z_{\parallel}(z) \sim q^N N^{\gamma_1} \eta_\parallel$ near the surface. A similar dependence on $z$ will be found for the number of configurations $Z_{\parallel}(z, z')$, when both ends of the chain are at the surface (or close to it), i.e., $Z_{\parallel}(z, z') \sim q^N N^{\gamma_1 \eta_\parallel}$. Another quantities of interest are $Z(0, z') \sim z^{-a} N^b \lambda_\parallel$, with $a_\parallel = \eta_\parallel - \eta_{\perp}$, $b_\parallel = -1 + \gamma_{11}$ for $z' \ll l N^\nu$ ($l$ is the effective monomer linear dimension, $l \to 0$ and $N \to \infty$) and the value of the mean number of polymer chain free ends in a layer between distances $z$ and $z + dz$ from the surface ($z = 0$) $Z(z) \sim z^{-a |N|^{b}}$ for $z \ll l N^\nu$, where the exponent $|a|$ characterizes density profiles for polymer chain ends and equals $|a| = (\gamma - \gamma_1)/\nu - \eta_\parallel/2$. The obtained results indicate that the presence of long-range correlated quenched disorder facilitates the process of adsorption of polymer chain on a hard wall. Besides, in a medium with long-range correlated quenched disorder the swelling of the polymer coil is governed by the exponent $\nu_{LR}$ that increases when the correlation of the disorder increases (i.e. $a$ decreases). When $a < a_{marg} < 2.3$, then a crossover to the collapse of the polymer is predicted. In the case of very strong correlation of the disorder we have to deal with a situation when such a polymer collapses and then adsorbs on a wall. The critical exponents which characterize the process of adsorption of long-flexible polymer chains inserted into the medium with long-range correlated disorder belong to the new universality class, different from that of the pure model.

Appendix 1

Applying the field-theoretical renormalization group (RG) approach we perform calculations in a double expansion in $\epsilon = 4 - d$ and in $\delta = 4 - a$ up to the one-loop order, as was proposed by Weinrib and Halperin\textsuperscript{16} for infinite systems. Thus, for the renormalization factor $Z_{\parallel}$ we obtain

$$Z_{\parallel} = 1 + \frac{\bar{v}}{4(1 + \epsilon)} + \frac{\bar{w}2^{-\delta} \sqrt{\pi}}{\cos(\frac{\pi}{2}(\delta - \epsilon))} \Gamma(\frac{\epsilon + \delta}{2}) \Gamma(\frac{\epsilon + \delta}{2}(1 + \delta)) + \frac{\delta - \epsilon}{2} \Gamma(\frac{\epsilon + \delta}{2} - 1) \Gamma(\frac{\epsilon + \delta}{2} - 2)(\frac{\Gamma(\frac{\epsilon + \delta}{2})}{\Gamma(\frac{\epsilon + \delta}{2} - 1)}) \right). \tag{A1.1}$$

Combining the renormalization factor $Z_{\parallel}$ together with the one-loop pieces of the beta functions, derived in\textsuperscript{7,17} we obtain for the surface critical exponent $\eta_{\parallel}^{sp}$

$$\eta_{\parallel}^{sp} = \frac{\bar{v}}{4(1 + \epsilon)} + \frac{\bar{w}}{2 \Gamma(\frac{\epsilon}{2})} \frac{\delta^{-\delta} \sqrt{\pi}}{\cos(\frac{\pi}{2}(\delta - \epsilon))} \Gamma(\frac{\epsilon + \delta}{2}) \Gamma(\frac{\epsilon + \delta}{2}(1 + \delta)) + \frac{\delta - \epsilon}{2} \Gamma(\frac{\epsilon + \delta}{2} - 1) \Gamma(\frac{\epsilon + \delta}{2} - 2)(\frac{\Gamma(\frac{\epsilon + \delta}{2})}{\Gamma(\frac{\epsilon + \delta}{2} - 1)}) \tag{A1.2}.$$
Appendix 2

The individual RG series expansions for the other critical exponents can be derived through standard surface scaling relations\textsuperscript{15} with $d = 3$

\begin{align}
\eta_\perp &= \frac{\eta + \eta_\parallel}{2}, \\
\beta_1 &= \nu \left( d - 2 + \eta_\parallel \right), \\
\gamma_{11} &= \nu \left( 1 - \eta_\parallel \right), \\
\gamma_1 &= \nu \left( 2 - \eta_\perp \right), \\
\Delta_1 &= \frac{\nu}{2} \left( d - \eta_\parallel \right), \\
\delta_1 &= \frac{\Delta_1}{\beta_1} = \frac{d + 2 - \eta}{d - 2 + \eta_\parallel}, \\
\delta_{11} &= \frac{\Delta_1}{\beta_1} = \frac{d - \eta_\parallel}{d - 2 + \eta_\parallel}.
\end{align}

Each of these critical exponents characterizes certain properties of the semi-infinite systems with long-range quenched disorder, in the vicinity of the critical point. The values $\nu$, $\eta$, and $\Delta = \nu(d + 2 - \eta)/2$ are the standard bulk exponents.

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