

## 6. Measurement of forces

### 6.1. Forces in air

#### 6.1.1. Meniscus force

Fig. 17 shows a force–displacement curve acquired on mica in air. Note that the hysteresis is very large. The jump-to-contact is rather small, while the jump-off-contact is quite large. Such a large adhesion is due to the meniscus force exerted by a thin layer of water vapor adsorbed on the sample surface, as schematically shown in Fig. 18. Such a layer barely affects attractive forces, whereas it prevents the tip from pulling off from the surface due to its high surface energy.

The Laplace pressure of the liquid meniscus is given by [15]:

$$P = \gamma_L \left( \frac{1}{r_1} + \frac{1}{r_2} \right) \cong \frac{\gamma_L}{r_{\text{eff}}}, \quad (6.1)$$

in which  $\gamma_L$  is the surface energy of the liquid and  $r_1$  and  $r_2$  are defined in Fig. 18. The Laplace pressure acts on the area [15]:

$$A \cong 2\pi R d \cong 2\pi R [r_{\text{eff}}(1 + \cos \theta) + u_M], \quad (6.2)$$

in which  $d$ ,  $\theta$ , and  $u_M$  are given in panel (a) of Fig. 18. The surfaces are pulled together with a force

$$F \cong 2\pi R \gamma_L \left( 1 + \cos \theta + \frac{u_M}{r_{\text{eff}}} \right) \quad (6.3)$$

Eq. (6.3) is the contribution of the Laplace pressure to the adhesion. Recently, Gao [96] has calculated formulas for different tip geometries, showing that the force is decreased by about 20% for a conical tip.

Mate et al. [97] have studied the meniscus force exerted by a perfluoropolyether polymer liquid film on silicon. Fig. 19 shows a typical force–displacement curve on such films.

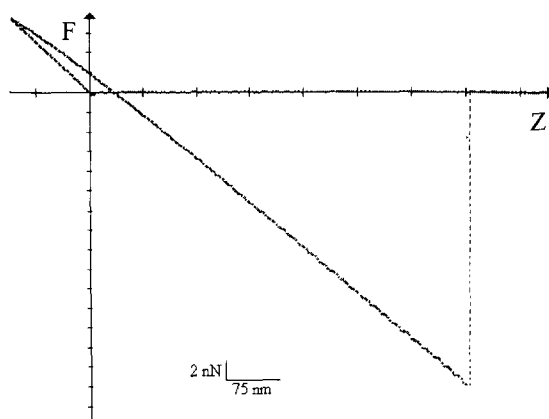


Fig. 17. Force–displacement curve on mica in air with a silicon nitride tip ( $R$  between 50 and 100 nm). The large jump-off-contact is due to the meniscus force. One can observe one of the most common artifacts in force–displacement curves, namely the inverse path effect (adapted from [128]).

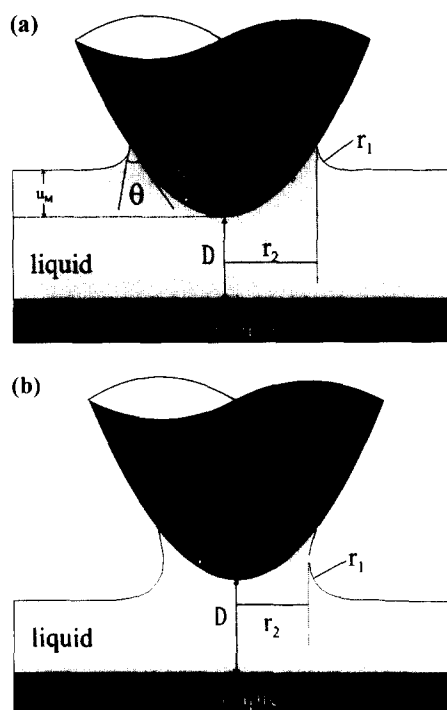


Fig. 18. The meniscus of the water vapor layer adsorbed on the sample surface onto a paraboloidal tip, during approach (panel (a)) and retraction (panel (b)).  $R$  is the radius of curvature of the tip,  $D$  is the distance from the sample,  $u_M$  is the penetration depth in the liquid layer,  $r_1$  and  $r_2$  the two radii that define the liquid meniscus,  $\theta$  the contact angle the liquid makes with the tip, and  $d$  the height of the meniscus relative to the end of the tip.

At great distances (right) the tip is far away from the surface and the deflection is zero. At 25 nm from the surface, while approaching the sample, the tip contacts the surface of the liquid film, that wicks up and forms a meniscus, resulting in an attractive force of 60 nN. After this first jump, the force on the tip is fairly constant, since the tip is penetrating through the liquid. When the tip contacts the

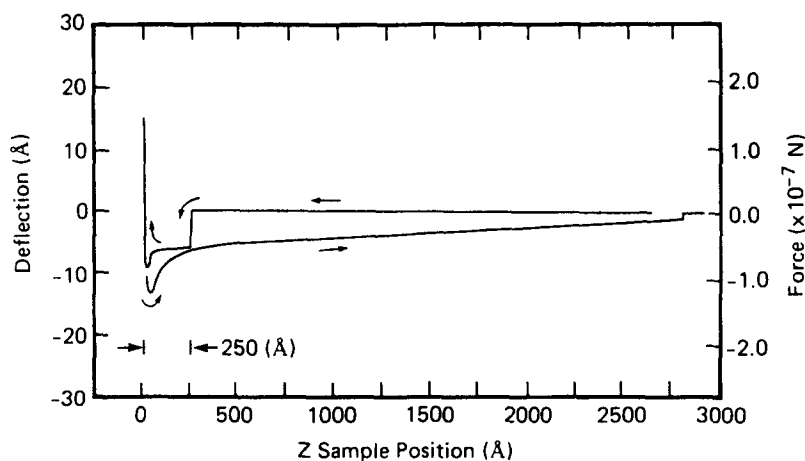


Fig. 19. Force-displacement curve on a perfluoropolyether polymer liquid film (reprinted with permission from [97]).

silicon, there is a second jump followed by the contact line. There are also two jumps-off-contact. The first jump, of 120 nN, occurs when the tip pulls off from silicon. The second, of  $\cong 10$  nN, when the tip detaches from the liquid film at  $\cong 280$  nm. In between, the force slowly decreases. The tip radius is approximately 230 nm. Force–distance curves allow to measure the effective meniscus radius  $r_{\text{eff}}$  and the thickness of the liquid film. Mate et al. have compared the thickness measured by the AFM with that measured by ellipsometry, showing that AFM always overestimates the film thickness, and that this is probably due to a thin film of liquid on the surface of the tip. Furthermore, the film thickness was acquired along a line of the sample in order to study the local liquid uniformity.

Thundat et al. [98] have observed the dependence of adhesion force on the relative humidity and have fitted the experimental data on meniscus force with the equation

$$F = - \frac{4\pi R \gamma_L \cos \theta}{[1 + D \log(H_r)] / 1.08 \cos \theta}, \quad (6.4)$$

in which  $D$  is the tip–sample distance in nm and  $H_r$  is the relative humidity. With a rather small tip ( $R = 15$  nm) the adhesion force is larger than 10 nN for  $H_r > 60\%$  and becomes smaller than 5 nN for  $H_r < 20\%$ . Later on Thundat et al. [99] reproduced the same fit for a clean and a contaminated tip, showing that adhesion forces are larger for the contaminated tip (70 nN at  $H_r > 50\%$  and 40 nN at  $H_r < 10\%$ ). When rinsing the contaminated tips in acetone, methanol or methylene chloride, a slight decrease ( $< 10\%$ ) in adhesion force is observed. Such a decrease is enhanced by cleaning the tips in UV ozone, in which case the adhesion reaches the values of the clean tips (see also Ref. [100]).

The increase of adhesion force with increasing relative humidity has been studied for silicon nitride tips on  $\text{SiO}_2$  substrates by Torii et al. [101].

Xu et al. [102] have studied the dependence of pull-off forces on relative humidity with Pt-covered tips on mica. The pull-off force is constant (2 nN) for  $H_r < 20\%$ , increases rapidly (8 nN) for  $H_r = 20\%$ , and then slowly decreases down to 3 nN for  $H_r = 95\%$ . They also verified that, depending on relative humidity, the pull-off force depends on contact time. At  $H_r = 42\%$ , the force goes from 87 nN for 1 ms contact time to 100 nN for a 10 s contact time. On the contrary, at  $H_r = 12\%$ , the pull-off force decreases with increasing contact time. Later on, Xu and Salmeron [103] have studied the effects of surface ions on the adhesion dependence on relative humidity.

Eastman and Zhu [104] have measured capillary forces between a mica substrate and surface modified tips, namely a  $\text{Si}_3\text{N}_4$  tip, a gold coated tip and a paraffin coated tip. The results show that the capillary force, and hence the total adhesion force diminishes when the tip is coated with a hydrophobic material. The value of the capillary force is 92 nN for the gold coated tip and  $-31$  nN for paraffin. The tip radius is about 100 nm. The authors do not obtain the same dependence of adhesion force as a function of relative humidity as Thundat et al., even for the  $\text{Si}_3\text{N}_4$  tip. They claim that this disagreement is due to the hydrophobicity of their mica substrate. Since the substrate is hydrophobic, the height of the water layer is limited and does not increase with relative humidity so that the adhesion force stays the same at all values of relative humidity.

The different dependence of meniscus force on relative humidity for hydrophilic and for hydrophobic substrates has been addressed by Fujihira et al. [105]. The authors study the adhesion force on a Si substrate partially covered with a hydrocarbon monolayer. They show that the adhesion force is about 5 nN on the hydrophilic Si substrate and 1 nN on the hydrophobic hydrocarbon substrate, using the same tip. Furthermore the meniscus force on silicon increases from 20 nN at 10% relative humidity to 40 nN at 60% relative humidity, while it is about 20 nN for any relative humidity value on hydrocarbon. This

latter measurement was performed using a tip with a larger curvature radius. Hence capillary force allows to distinguish hydrophobic and hydrophilic materials. A similar set of measurements has been performed by Binggeli and Mate [106] on silicon, carbon and perfluoropolyether and by Olsson et al. [107] between a tungsten tip and a methylated  $\text{SiO}_2$  surface. In this latter work, dynamic force–distance curves also are acquired, showing that, when the meniscus forms and the capillary force onsets, the amplitude of the oscillations abruptly goes from the free amplitude to an almost zero amplitude via a sharp and strong discontinuity. Dynamic force–distance curves are, therefore, more suitable for the detection of the distance of formation of menisci. This capability of dynamic force–displacement curves has been exploited by Friedenbergs and Mate [108] in order to study the meniscus force exerted by different layers of polydimethylsiloxane (PDMS). Dynamic force–displacement curves are acquired varying the penetration distances in the PDMS layer and the thickness of the PDMS layer. The experimental results are in rather good agreement with theoretical predictions. Colchero et al. [109] have acquired force–displacement curves in the static and in the dynamic modes simultaneously. They have shown that static force–displacement curves reveal the liquid–solid interface, while the dynamic force–distance curve is able to reveal the presence of the liquid–vapor interface.

The study of capillary forces, of great importance also for the comprehension of the mechanisms involved in AFM imaging, has revealed the measurement of their effect as a useful technique for distinguishing materials with different hydrophobicity (see also Section 7). The different techniques reviewed here may be used effectively for the study of wetting properties of lubricants.

Water meniscus force exceeds all other forces, and in particular, it masks the Van der Waals force. Inserting  $\gamma_L = 72 \text{ mJ m}^{-2}$ ,  $R = 30 \text{ nm}$  and  $d = 2 \text{ nm}$  into Eq. (6.3) the meniscus force is approximately 30 nN, whereas the adhesion force due to Van der Waals attraction for normal materials is typically only 20 nN. In order to measure Van der Waals force and/or other surface forces, the meniscus force must be eliminated, either by removing the water layer by working in a low humidity environment (such as dry nitrogen or vacuum), or by dipping both sample and tip in a liquid environment [57].

### 6.1.2. Coulomb force

The measure of Coulomb forces until now has been of little interest, but it can be employed in order to study the tip shape. Hao et al. [21] have studied Coulomb forces by modeling the tip–sample system as a sphere on a flat surface and as a sphere-ended conical tip on a flat sample, as shown in Fig. 20. In the first case, the force is given by

$$F_0^{\text{sp}} = \pi \varepsilon_0 V^2 \frac{R}{D} \quad (6.5a)$$

for  $R/D \gg 1$  and by

$$F_\infty^{\text{sp}} = \pi \varepsilon_0 V^2 \left( \frac{R}{D} \right)^2 \quad (6.5b)$$

for  $R/D \ll 1$ .  $V$  is the voltage difference between the tip and sample,  $R$  the radius of the sphere and  $D$  is the tip–sample distance.

In the case of a sphere-ended cone on a flat surface, the force may be calculated by replacing the equipotential conducting surfaces with their equivalent image charges. For small aperture angles

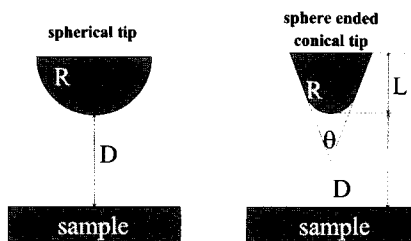


Fig. 20. The two tip–sample systems employed by Hao et al. [21] for the calculation of Coulomb force.  $R$  is the radius of the sphere,  $D$  the tip–sample distance,  $L$  the length of the sphere-ended conical tip, and  $\theta$  the aperture angle.

( $\theta \leq \pi/9$ ) the cone may be approximated by a charged line of constant charge density  $\lambda_0$  given by

$$\lambda_0 = 4\pi\epsilon_0 V \left[ \ln \left( \frac{1 + \cos \theta}{1 - \cos \theta} \right) \right]^{-1}. \quad (6.6)$$

The resulting force is given by

$$F^c \cong \frac{\lambda_0^2}{4\pi\epsilon_0} \ln \left( \frac{L}{4D} \right), \quad (6.7)$$

in which  $L$  is the cone length ( $L \ll D \ll R$ ).

Hao et al. have measured the Coulomb forces between a tungsten tip and graphite. By fitting the experimental data, the authors have been able to determine the curvature radius of different tips (270 and 27.5 nm, in good agreement with scanning electron microscope measurements). Furthermore, they have shown that Eq. (6.5a) fits well the measured forces at small tip–sample separations, whereas at larger distances Eq. (6.7) is needed (see Fig. 21). The disagreement between experimental data and Eq. (6.5b) increases with decreasing tip dimensions.

Burnham et al. [110] have studied another kind of Coulomb-like force which arises from the patch charges distribution on the tip and sample, i.e., from regions of different surface charge density

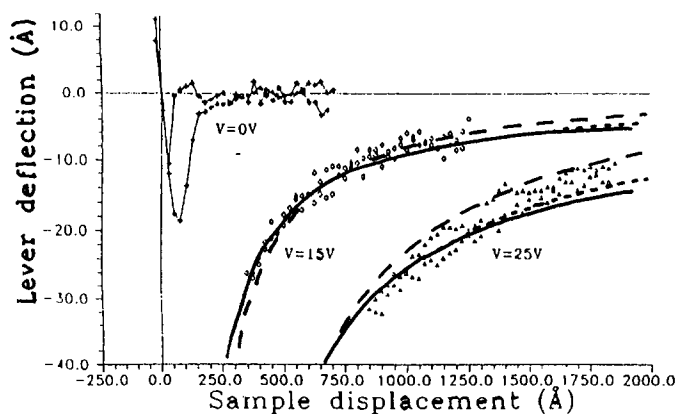


Fig. 21. Force–displacement curves on graphite. Coulomb forces at  $V=0$  V,  $V=15$  V and  $V=25$  V are displayed. The theoretical results for the sphere-ended conical tip (continuous line), the spherical tip (— — —) and the simple  $R/D$  approximation (- - -) are also shown (reprinted with permission from [21]; copyright 1991 American Vacuum Society).

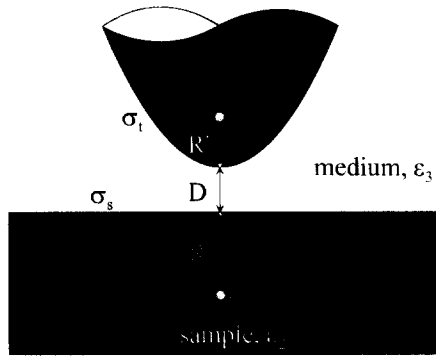


Fig. 22. The tip-sample system employed by Burnham et al. [110] for the calculation of patch charges force. The tip and sample surface charges are represented by two image charges  $Q_t$  and  $Q_s$  placed at the distances  $R'$  and  $R''$ .  $D$  is the tip-sample distance.  $\varepsilon_1$ ,  $\varepsilon_2$  and  $\varepsilon_3$  are the dielectric constants of the tip, the sample, and the medium, respectively.

interacting via a long range force law. The patch charges distribution can be modeled by means of the image charge method, as shown in Fig. 22. Consider a spherical tip of radius  $R$  on a flat sample with surface charge densities  $\sigma_t$  and  $\sigma_s$  represented by an image charge  $Q_t$  at a distance  $R'$  inside the tip and by an image charge  $Q_s$  at a distance  $R''$  inside the sample. If  $R \gg R' \gg D$  and  $R'' \gg D, R, R'$ , then the force is given by

$$4\pi\varepsilon_0 F = -\frac{Q_t}{4R'^2} \left(1 - \frac{2D}{R'}\right) \left(\frac{\varepsilon_2 - \varepsilon_3}{\varepsilon_2 + \varepsilon_3}\right) + \frac{Q_t Q_s}{RR''} \left(1 - \frac{2R'}{R} - \frac{4D}{R}\right) \left(\frac{\varepsilon_1 - \varepsilon_3}{\varepsilon_1 + \varepsilon_3}\right) \left(\frac{\varepsilon_2 - \varepsilon_3}{\varepsilon_2 + \varepsilon_3}\right), \quad (6.8)$$

where  $\varepsilon_1$  and  $\varepsilon_2$  are the dielectric constants of the tip and sample and  $\varepsilon_3$  is the dielectric constant of the medium.

Patch charge densities are negligible compared to double-layer charge densities (see Section 6.3.1). The patch charge force is larger than the Van der Waals force and has the same dependence on the dielectric constants. Note that, if the geometrical constants  $R$  and  $R'$  are much larger than  $D$ , Eq. (6.8) becomes effectively independent on the distance. There is evidence that this occurs when the ratio of the probe radius to the distance is  $10^3$ .

This kind of force has been measured by Burnham et al. [110] and by Agraït et al. [59]. This latter experiment has been performed in vacuum, at liquid helium temperature (4.2 K), using a gold tip on a gold sample. The presence of contaminants can therefore be excluded, allowing the different forces to be distinguished from each other.

## 6.2. Van der Waals force

### 6.2.1. Theory

In the theoretical subsections only the issues necessary to the comprehension of the experimental results are reviewed. The discussion is based on the derivations given in [15, 111–113].

The Van der Waals force between atoms and/or molecules is the sum of three different forces, all proportional to  $1/r^6$  ( $r$  is the distance between the atoms or molecules). These forces are: (a) the orientation or Keesom force, (b) the induction or Debye force, and (c) the dispersion or London force.

- (a) The Keesom force is the angle-averaged dipole–dipole interaction between two atoms or molecules and its potential is [15]:

$$w_K(r) = -\frac{u_1^2 u_2^2}{3(4\pi\epsilon_0\epsilon)^2 k_B T r^6} = -\frac{C_K}{r^6}, \quad (6.9)$$

where  $u_1$  and  $u_2$  are the dipole moments of the molecules,  $\epsilon$  the dielectric constant of the medium,  $k_B$  the Boltzmann constant and  $T$  is the temperature. Eq. (6.9) holds for  $k_B T > u_1 u_2 / 4\pi\epsilon_0\epsilon r^3$  [15].

- (b) The Debye force is the angle-averaged dipole-induced dipole interaction between two atoms or molecules:

$$w_D(r) = -\frac{u_1^2 \alpha_{02} + u_2^2 \alpha_{01}}{(4\pi\epsilon_0\epsilon)^2 r^6} = -\frac{C_D}{r^6} \quad (6.10)$$

in which  $\alpha_{01}$  and  $\alpha_{02}$  are the electronic polarizabilities of the molecules. Eq. (6.9) may be obtained from Eq. (6.10) by replacing the electronic polarizability  $\alpha_0$  with the orientation polarizability  $\alpha = u^2 / 3k_B T$ .

- (c) The dispersion force is the most important contribution to Van der Waals force, because it acts between all molecules or atoms, on a distance that ranges from more than 10 nm down to 2 Å. It may be attractive or repulsive. The dispersion force is the instantaneous dipole-induced dipole interaction, and is of quantum-mechanical nature. An expression of this force has been calculated in 1930 by London [114]. The corresponding potential is

$$w_L(r) = -\frac{3}{2} \frac{\alpha_{01} \alpha_{02}}{(4\pi\epsilon_0)^2 r^6} \frac{(h\nu_1)(h\nu_2)}{h\nu_1 + h\nu_2} = -\frac{C_L}{r^6}, \quad (6.11)$$

in which  $h\nu_1$  and  $h\nu_2$  are the first ionization potentials of the molecules and  $h$  is the Planck constant.

Thus the total Van der Waals potential can be written as

$$w_{vdW}(r) = -\frac{C_K + C_D + C_L}{r^6} = -\frac{1}{(4\pi\epsilon_0)^2 r^6} \left[ (u_1^2 \alpha_{02} + u_2^2 \alpha_{01}) + \frac{u_1^2 u_2^2}{3k_B T} + \frac{3\alpha_{01} \alpha_{02} (h\nu_1)(h\nu_2)}{2(h\nu_1 + h\nu_2)} \right]. \quad (6.12)$$

In 1963 McLachlan [115] presented a generalized theory of the Van der Waals force between infinite media 1 and 2 separated by a medium 3 in the form:

$$w_{vdW}(r) = -\frac{6k_B T}{(4\pi\epsilon_0)^2 r^6} \sum_{n=0}^{\infty} \frac{\alpha_1(i\nu_n) \alpha_2(i\nu_n)}{\epsilon_3^2(i\nu_n)}, \quad (6.13)$$

in which the 0-term of the sum has to be multiplied by 1/2,  $\alpha_1(i\nu_n)$  and  $\alpha_2(i\nu_n)$  are the polarizabilities of molecules, and  $\epsilon_3(i\nu_n)$  is the dielectric constant of the medium at the frequencies  $i\nu_n = 2\pi i n k_B T / h$ . The total polarizability in vacuum at the frequency  $i\nu_n$  is [116]:

$$\alpha(i\nu_n) = \frac{\alpha_0}{1 + (\nu_n/\nu_1)^2} + \frac{u^2}{3k_B T(1 + \nu_n/\nu_{rot})}, \quad (6.14)$$

in which  $\nu_1$  is the ionization frequency and  $\nu_{rot}$  is the average rotational relaxation frequency.

The first term ( $n = 0$ ) of Eq. (6.13) is:

$$(w_{\text{vdW}}(r))_{v=0} = -\frac{3k_{\text{B}}T}{(4\pi\epsilon_0)^2 r^6} \alpha_1(0)\alpha_2(0) = -\frac{3k_{\text{B}}T}{(4\pi\epsilon_0)^2 r^6} \left( \frac{u_1^2}{3k_{\text{B}}T} + \alpha_{01} \right) \left( \frac{u_2^2}{3k_{\text{B}}T} + \alpha_{02} \right). \quad (6.15)$$

It contains the Keesom and Debye energies. This term acts only between polar molecules and is called polar or entropic contribution.

The term at non-zero frequencies is called dispersion contribution. This sum entails the calculation of  $\alpha(i\nu_n)$  at the discrete frequencies  $\nu_n$  and hence the knowledge of the absorption spectrum of the materials over the entire real frequency range,  $0 \leq \nu < \infty$ . With a few exceptions, only partial data are available, and some approximate methods have been developed, such as the Ninham and Parsegian method [117]. In this method  $\alpha(i\nu_n)$  is considered as determined only by electronic polarizability, since, at normal temperatures, the first frequency  $\nu_{n=1} \cong 4 \times 10^{13}$  Hz is already much greater than  $\nu_{\text{rot}} \cong 10^{11}$  Hz. Further, since all the frequencies  $\nu_n$  are very close in the UV region, the sum in Eq. (6.13) can be replaced [15] by an integration over  $dn = (h/2\pi k_{\text{B}}T) d\nu$ :

$$(w_{\text{vdW}}(r))_{\nu>0} = -\frac{3h}{(4\pi\epsilon_0)^2 \pi r^6} \int_{\nu_1}^{\infty} \alpha_1(i\nu)\alpha_2(i\nu) d\nu \cong -\frac{3}{2} \frac{\alpha_{01}\alpha_{02}}{(4\pi\epsilon_0)^2 r^6} \frac{h\nu_1\nu_2}{\nu_1 + \nu_2}, \quad (6.16)$$

yielding Eq. (6.11).

McLachlan theory allows to calculate dispersion forces also for atoms or molecules of dielectric constant  $\epsilon_1$  in a medium of dielectric constant  $\epsilon_3$ . In this case it is necessary to consider the excess polarizability of the molecule of radius  $R_{\text{m}}$  given by the Clausius–Mossotti equation [118]:

$$\alpha_{\text{exc}}(\nu) = 4\pi\epsilon_0\epsilon_3(\nu) \frac{\epsilon_1(\nu) - \epsilon_3(\nu)}{\epsilon_1(\nu) + 2\epsilon_3(\nu)} R_{\text{m}}^3. \quad (6.17)$$

Inserting Eq. (6.17) into Eqs. (6.15) and (6.16) the total Van der Waals interaction, Eq. (6.13), becomes the sum of two terms given by

$$(w_{\text{vdW}}(r))_{v=0} = -\frac{3k_{\text{B}}TR_{\text{m1}}^3R_{\text{m2}}^3}{r^6} \left( \frac{\epsilon_1(0) - \epsilon_3(0)}{\epsilon_1(0) + 2\epsilon_3(0)} \right) \left( \frac{\epsilon_2(0) - \epsilon_3(0)}{\epsilon_2(0) + 2\epsilon_3(0)} \right), \quad (6.18a)$$

and

$$(w_{\text{vdW}}(r))_{v>0} = -\frac{3hR_{\text{m1}}^3R_{\text{m2}}^3}{\pi r^6} \int_0^{\infty} \frac{\epsilon_1(i\nu) - \epsilon_3(i\nu)}{\epsilon_1(i\nu) + 2\epsilon_3(i\nu)} \frac{\epsilon_2(i\nu) - \epsilon_3(i\nu)}{\epsilon_2(i\nu) + 2\epsilon_3(i\nu)} d\nu. \quad (6.18b)$$

If the dielectric medium has one strong absorption peak at the frequency  $\nu_e$  (that is usually different from the frequency  $\nu_1$  of the isolated molecule),  $\epsilon(\nu)$  can be written as

$$\epsilon(\nu) = 1 + \frac{n^2 - 1}{1 + (\nu/\nu_e)^2}, \quad (6.19)$$

in which  $n$  is the refractive index. Substituting in Eq. (6.18b) we obtain:

$$(w_{\text{vdW}}(r))_{v>0} = -\frac{\sqrt{3}h\nu_e R_{\text{m1}}^3 R_{\text{m2}}^3}{2 r^6} \frac{(n_1^2 - n_3^2)(n_2^2 - n_3^2)}{\sqrt{(n_1^2 + 2n_3^2)}\sqrt{(n_2^2 + 2n_3^2)}[\sqrt{(n_1^2 + 2n_3^2)} + \sqrt{(n_2^2 + 2n_3^2)}]}. \quad (6.20)$$



Hence the total Van der Waals force between two identical molecules is:

$$w_{\text{vdw}}(r) = - \left[ 3k_{\text{B}}T \left( \frac{\varepsilon_1(0) - \varepsilon_3(0)}{\varepsilon_1(0) + 2\varepsilon_3(0)} \right)^2 + \frac{\sqrt{3}}{4} h\nu_e \frac{(n_1^2 - n_3^2)^2}{\sqrt{(n_1^2 + 2n_3^2)^3}} \frac{R_{\text{ml}}^3}{r^6} \right]. \quad (6.21)$$

From the above equations, some conclusions about Van der Waals force can be drawn.

- (1) Whenever  $h\nu_e \gg k_{\text{B}}T$ , as in vacuum, the dispersion term is larger than the polar term.
- (2) In a medium, the Van der Waals force is greatly reduced.
- (3) The Van der Waals dispersion (polar) force in a medium may be attractive, vanishing or repulsive. It is repulsive whenever  $n_3(\varepsilon_3)$  is intermediate between  $n_1(\varepsilon_1)$  and  $n_2(\varepsilon_2)$ . It vanishes if  $n_3(\varepsilon_3)$  matches  $n_1(\varepsilon_1)$  or  $n_2(\varepsilon_2)$ . The force between two identical molecules is always attractive.

Some other features of Van der Waals force are noteworthy.

(1) The Van der Waals force is anisotropic. As a matter of fact, the polarizabilities of the majority of molecules are anisotropic, i.e., they have different values for different molecular directions. One of the consequences of the anisotropy of polarizabilities is that the dispersion force between two molecules depends on their mutual orientation. This effect is important in solids and liquid crystals, where the thermal motion is reduced.

(2) The Van der Waals force is non-additive. The force between two molecules is affected by other molecules nearby, since they behave like a medium.

(3) The Van der Waals force is affected by retardation effects. This effect is described in detail in [119] and only a qualitative description is given here. When two atoms are far away, the time taken by the field of the first atom to reach the second one and come back may be comparable to the period of the fluctuating dipole. In this case, the reflected field finds that the orientation of the first dipole is now different. At large separation, because of the retardation effect, the power law index increases by unity. Thus, with increasing separation, the dispersion force begins to decay as  $-1/r^7$  at distances larger than 100 nm. Such retardation effects turn out to be very important in liquids, where the speed of light is low, and the effect appears at distances smaller than 5 nm. Only the dispersion term is retarded, whereas the polar one is non-retarded at all distances. In the presence of a medium, the Van der Waals force has three distance regimes. At small distances the force law is  $1/r^6$ . At intermediate distances, where the retarded dispersion dominates the Van der Waals force, the force law is  $1/r^7$ . At larger distances, the polar term, that is always non-retarded, turns out to be dominant and the force law is again  $1/r^6$ .

In order to model the interactions taking place in AFM, it is necessary to study interactions between surfaces, rather than between atoms or molecules. Let us begin with a pair potential in the form  $w(r) = -C/r^n$ . Assuming that the potential is additive, the interaction energies between macroscopic bodies may be obtained via integration. In the case of two interacting spheres of radius  $R_1$  and  $R_2$  at the distance  $D$ , the force  $F(D)$  can be obtained by integrating over small circular sections of surface  $2\pi x dx$  on both spheres [120] (Fig. 23):

$$F(D) = \int_D^\infty 2\pi x dx f(Z), \quad (6.22)$$

in which  $f(Z)$  is the normal force per unit area. From the Chord theorem  $x^2 \cong 2R_1z_1 = 2R_2z_2$  and

$$Z = D + z_1 + z_2 = D + \frac{x^2}{2} \left( \frac{1}{R_1} + \frac{1}{R_2} \right), \quad (6.23)$$

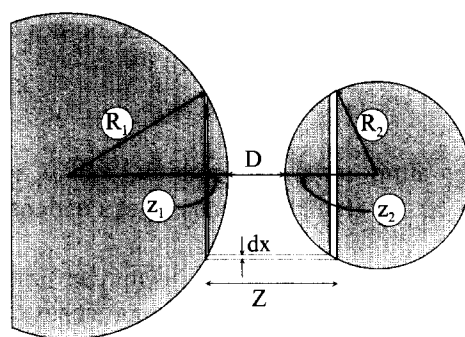


Fig. 23. Geometry employed in the Derjaguin approximation. The two spheres of radii  $R_1$  and  $R_2$  are at a distance  $D$ . The force is integrated over small circular sections of radius  $x$  and height  $dx$  at distances  $Z = D + z_1 + z_2$ .

and hence

$$dZ = \left( \frac{1}{R_1} + \frac{1}{R_2} \right) x dx. \quad (6.24)$$

Finally, the force is given by

$$F(D) \cong \int_D^\infty 2\pi \frac{R_1 R_2}{R_2 + R_2} f(Z) dZ = 2\pi \frac{R_1 R_2}{R_1 + R_2} W(D), \quad (6.25)$$

in which  $W(D)$  is the interaction potential between two flat surfaces.

Eq. (6.25) is known as the Derjaguin approximation, and holds whenever the interaction range and the separation  $D$  are much smaller than  $R_1$  and  $R_2$ . It should be emphasized that in Eq. (6.25) it is not necessary to specify explicitly the type of interaction  $f(Z)$ . This fact implies that Eq. (6.25) is correct for any kind of force, whether attractive, repulsive, or oscillating. From the Derjaguin approximation, by putting  $R_2 \gg R_1$ , it is possible to obtain the force between a sphere and a flat surface:

$$F(D) = 2\pi R_1 W(D). \quad (6.26)$$

Van der Waals energies between macroscopic bodies in vacuum may be computed via integration only in the approximation that the Van der Waals force is considered additive and non-retarded. The interaction laws obtained via integration are listed in Table 1 for the most common geometries.

Table 1

Van der Waals interactions for the most common geometries.  $R$  is the radius of spheres,  $D$  the distance between the interacting bodies, and  $A$  the Hamaker constant.

Geometry	Interaction law
Two spheres	$W = -\frac{A}{6D} \frac{R_1 R_2}{R_1 + R_2}$
Sphere-flat surface	$W = -\frac{AR}{6D}$
Two flat surfaces	$W = -\frac{A}{12\pi D^2}$ per unit area

All interaction laws depend on geometrical features of the bodies and on the Hamaker constant  $A$ , which includes all physico-chemical information [121]:

$$A = \pi^2 C \rho_1 \rho_2, \quad (6.27)$$

in which  $C$  is the constant in the atom–atom pair potential and  $\rho_1$  and  $\rho_2$  are the number of atoms per unit volume. Typical values of the Hamaker constant of condensed phases in vacuum are about  $10^{-19}$  J.

In the following, we will show how the Hamaker constant can be calculated. However, in order to determine Hamaker constants from measured curves, two approaches can be followed. The first method consists in fitting one of the force laws listed in Table 1 to the attractive part of the force–distance curve. This procedure can be somewhat difficult due to the short interaction range. The second method adapts Eq. (1.9) to the Van der Waals force and to the appropriate geometry. Thus, the Hamaker constant can be determined from the jump-to-contact force.

For  $n = 2$  and  $C = AR/6$ , i.e., for a sphere on a flat surface, Eq. (1.9) becomes

$$\begin{aligned} (\delta_c)_{\text{jtc}} &= \frac{1}{2} \sqrt[3]{\frac{AR}{3\beta^2 k_c}} \\ D_{\text{jtc}} &= 2\beta(\delta_c)_{\text{jtc}} = \sqrt[3]{\frac{\beta AR}{3k_c}}. \end{aligned} \quad (6.28)$$

Eq. (6.28) permits to calculate  $A$  and  $R$  from the measured  $(\delta_c)_{\text{jtc}}$  and  $D_{\text{jtc}}$  once  $\beta$  and  $k_c$  are known.

By assuming the additivity of the Van der Waals force, the influence of nearby atoms on the couple of interacting atoms is neglected, thus introducing large errors in the calculation of the Hamaker constant. To circumvent this problem, Lifshitz [122] presented an alternative approach in which each body is treated as a continuum with certain dielectric properties. This approach automatically incorporates many-body effects which are neglected in the microscopic approach. All the expressions in Table 1 remain valid. Only the computation of the Hamaker constant changes.

In Lifshitz theory, the Hamaker constant for the interaction of media 1 and 2 across the medium 3 may be expressed in terms of McLachlan equation (6.13) as

$$\begin{aligned} A &= \pi^2 C \rho_1 \rho_2 = \frac{6\pi^2 k_B T \rho_1 \rho_2}{(4\pi\epsilon_0)^2} \sum_0^\infty \frac{\alpha_1(i\nu_n) \alpha_2(i\nu_n)}{\epsilon_3^2(i\nu_n)} \\ &= \frac{3}{2} k_B T \sum_0^\infty \left[ \frac{\epsilon_1(i\nu_n) - \epsilon_3(i\nu_n)}{\epsilon_1(i\nu_n) + \epsilon_3(i\nu_n)} \right] \left[ \frac{\epsilon_2(i\nu_n) - \epsilon_3(i\nu_n)}{\epsilon_2(i\nu_n) + \epsilon_3(i\nu_n)} \right] \end{aligned} \quad (6.29)$$

Substituting the sum with an integration, as already done in Eq. (6.16), yields

$$A \cong \frac{3}{4} k_B T \frac{\epsilon_1 - \epsilon_3}{\epsilon_1 + \epsilon_3} \frac{\epsilon_2 - \epsilon_3}{\epsilon_2 + \epsilon_3} + \frac{3h}{4\pi} \int_{\nu_1}^\infty \frac{\epsilon_1(i\nu_n) - \epsilon_3(i\nu_n)}{\epsilon_1(i\nu_n) + \epsilon_3(i\nu_n)} \frac{\epsilon_2(i\nu_n) - \epsilon_3(i\nu_n)}{\epsilon_2(i\nu_n) + \epsilon_3(i\nu_n)} d\nu. \quad (6.30)$$

The first term is the zero frequency contribution and includes Debye and Keesom forces. The second term is the dispersion contribution.

In order to compute Hamaker constants it is necessary to know the dependence of the dielectric constant on the frequency. Hough and White [113], following Ninham and Parsegian [117], have

approximated  $\varepsilon(i\nu)$  in the form:

$$\varepsilon(i\nu) = 1 + \frac{C_{\text{IR}}}{1 + \nu/\nu_{\text{rot}}} + \frac{C_{\text{UV}}}{1 + \nu^2/\nu_e^2}, \quad C_{\text{IR}} \cong \varepsilon_0 - n^2; \text{ and } C_{\text{UV}} \cong n^2 - 1. \quad (6.31)$$

Such an approximation holds for materials with sharp, non-overlapping absorption peaks. The infrared term can be neglected when  $\varepsilon_0 \cong n^2$ . Using the Hough and White approximation, and neglecting the infrared term, Eq. (6.30) becomes

$$\begin{aligned} A_{\text{TOT}} &= A_{\nu=0} + A_{\nu>0} \\ &\cong \frac{3}{4} k_B T \frac{\varepsilon_1 - \varepsilon_3}{\varepsilon_1 + \varepsilon_3} \frac{\varepsilon_2 - \varepsilon_3}{\varepsilon_2 + \varepsilon_3} + \frac{3h\nu_e}{8\sqrt{2}} \frac{(n_1^2 - n_3^2)(n_2^2 - n_3^2)}{\sqrt{n_1^2 + n_3^2} \sqrt{n_2^2 + n_3^2} [\sqrt{n_1^2 + n_3^2} + \sqrt{n_2^2 + n_3^2}]} \end{aligned} \quad (6.32)$$

For two identical media ( $\varepsilon_1 = \varepsilon_2 \neq \varepsilon_3, n_1 = n_2 \neq n_3$ ), Eq. (6.32) becomes

$$A \cong \frac{3}{4} k_B T \left( \frac{\varepsilon_1 - \varepsilon_3}{\varepsilon_1 + \varepsilon_3} \right)^2 + \frac{3h\nu_e}{16\sqrt{2}} \frac{(n_1^2 - n_3^2)^2}{\sqrt{(n_1^2 + n_3^2)^3}}. \quad (6.33)$$

Eq. (6.33) exhibits three interesting aspects.

1. The Van der Waals force between two identical bodies in a medium is always attractive ( $A$  is positive), whereas the force between two different bodies may be attractive or repulsive.

2. The Van der Waals force between any two condensed bodies in vacuum or in air ( $\varepsilon_3 = 1$  and  $n_3 = 1$ ) is always attractive.

3. The polar term cannot be larger than  $(3/4)k_B T$ , which is about  $3.6 \times 10^{-21}$  J at  $T = 300$  K.

Bergström [112] has evaluated the accuracy of Ninham–Parsegian and Hough–White approximations, and the importance of the infrared contribution.

For interactions between conductive bodies such as metals, Eq. (6.33) cannot be applied, since  $\varepsilon$  is infinite. For two metals in vacuum the Hamaker constant is [15]:

$$A \cong \frac{3}{8\sqrt{2}} h \frac{\nu_{e1}\nu_{e2}}{\nu_{e1} + \nu_{e2}} \cong 4 \cdot 10^{-19} \text{ J}, \quad (6.34)$$

in which  $\nu_{e1}$  and  $\nu_{e2}$  are the plasma frequencies of the two metals.

At large separations, the retardation effect becomes important, and a retarded Hamaker constant  $A_r$  has to be considered [111, 123]. The entropic Hamaker constant scales as  $k_B T$  and the non-retarded Hamaker constant scales as  $h\nu_e$ . Similarly, the retarded Hamaker constant scales as  $hc$  ( $c$  is the speed of light). The characteristic wavelength

$$\lambda_{132} = 6\pi A_r / A \quad (6.35)$$

indicates the onset of retardation.  $A_r$  depends on low-frequency indices of refraction  $n^*$ . It is negative for  $n_1^* < n_3^* < n_2^*$  and positive for  $n_3^* < n_1^*$  or  $n_3^* > n_2^*$ . Anyway, since there is no correlation between  $n^*$  and  $n$ , the Van der Waals force may be attractive at small distances and repulsive at great distances, and vice versa.

Hamaker constants for materials commonly used in AFM are listed in Table 2. Usual AFM tips are made out of silicon nitride or silicon, one common substrate is mica and the colloidal probes are usually silica spheres. A complete list of Hamaker constants calculated with the full Hough and White

Table 2

Hamaker constants ( $\times 10^{-20}$  J) through air (*italic values*) or water (**bold values**) for most common AFM materials as calculated by (a) Bergström [112] using the Hough and White expression and the Prieve-Russel representation for water, (b) Bergström using the approximate expressions Eqs. (6.31) or (6.32), (c) Senden and Drummond [124] using Hough and White method and Gingell-Parsegian representation of water.

	Si <sub>3</sub> N <sub>4</sub>	silicon	mica	silica
Si <sub>3</sub> N <sub>4</sub>	<i>16.7<sup>(a)</sup>, 16.2<sup>(b)</sup>, <b>4.85<sup>(a)</sup>, 6.07<sup>(c)</sup></b></i>	<i>16.84<sup>(c)</sup>, <b>6.75<sup>(c)</sup></b></i>	<i>12.8<sup>(a)</sup>, <b>2.45<sup>(a)</sup>, 3.40<sup>(c)</sup></b></i>	<i>10.8<sup>(a)</sup>, <b>1.17<sup>(a)</sup>, 1.9<sup>(c)</sup></b></i>
silicon		<i>18.65<sup>(c)</sup>, <b>9.75<sup>(c)</sup></b></i>	<i>12.63<sup>(c)</sup>, <b>3.48<sup>(c)</sup></b></i>	<i>10.26<sup>(c)</sup>, <b>1.92<sup>(c)</sup></b></i>
mica			<i>9.86<sup>(a)</sup>, 9.64<sup>(b)</sup>, <b>1.34<sup>(a)</sup>, 2.01<sup>(c)</sup></b></i>	<i>8.01<sup>(a)</sup>, <b>0.69<sup>(a)</sup>, 1.19<sup>(c)</sup></b></i>
silica				<i>6.5<sup>(a)</sup>, 6.39<sup>(b)</sup>, <b>0.46<sup>(a)</sup>, 0.77<sup>(c)</sup></b></i>

approximation is given by Bergström [112]. In Table 2 the values calculated by Senden and Drummond [124] are also included. In the case of water as a medium, Hamaker constants may differ up to 67% due to the different models of water used in the calculations. The exclusion of the infrared term brings an error of about 2%.

### 6.2.2. Experimental results

At the very beginning of the study of tip-sample interactions with the AFM, the importance of capillary force had not been realized, and some researchers pretended to measure Van der Waals force and sample surface energies in air. Weisenhorn et al. [57] were the first to realize the necessity of working in water if one would like to perform measurement in absence of capillary force.

Nevertheless, a lot of measurements have been performed in dry nitrogen. There are some common drawbacks affecting these initial experiments.

1. The tips are rather large, so that the interaction is averaged over a large area and details are lost. Furthermore, any asperity on the tip turns to determine the interaction.

2. Surface contaminants are likely to store up on a very large tip and cannot be easily eliminated.

3. Maugis theory was still unknown, so that the elastic contact is modeled according to the JKR theory or according to the DMT analysis.

The previous considerations do not hold for the study of Eastman and Zhu [104] and for that of Hutter and Bechhoefer [95]. The work of Eastman and Zhu [104] is focused on determining separately the contributions of capillary and Van der Waals forces (see Section 6.1). The work of Hutter and Bechhoefer [95] is the first measurement of the transition from the non-retarded regime to the retarded regime of the Van der Waals force. The power law is  $-2$  ( $-2.19$  is the experimental value) for distances smaller than 16 nm and increases to  $-3$  ( $-2.92$ ) for distances larger than 20 nm. The exponent  $-2$  is characteristic for the interaction between a sphere and a plane. The increased exponent at larger distances is due to the retardation effect. Table 3 contains a list of measurements performed in air or in dry nitrogen. Even if measured quantities have a large uncertainty, they often show a valid trend with surface energies.

Fig. 24 shows a curve acquired on mica in water. Both attractive and adhesion forces become nearly 10 times smaller than in air and the Van der Waals force makes the greatest contribution to adhesion since the meniscus has been removed. Also the pull-off distance is reduced by a factor of 10.

Table 4 provides a list of measurements in water or water solutions.

Table 3

Measured attractive force  $F_{\text{attr}}$ , adhesion force  $F_{\text{ad}}$ , and surface energy  $\gamma_s$  for experiments in air or in dry nitrogen: (a) The tip material and dimensions are unknown in this work, (b) The tips are colloidal particles, (c) PTFE = poly-tetrafluoroethylene.

Ref.	Tip radius (nm)	Tip material/Medium/Sample	Measured (calculated)		
			$F_{\text{attr}}$ (nN)	$\gamma_s$ (mJ/m <sup>2</sup> )	$F_{\text{ad}}$ (nN)
[18]	270	SiO <sub>2</sub> /air/graphite, LiF			100, 160
[57]	(a)	(a)/air/mica			400
[44]	50–100	tungsten/dry N <sub>2</sub> /silicon			50
[42]	100–200	tungsten/dry N <sub>2</sub> /gold	$\cong 750$	270–540	$\cong 750$
		tungsten/dry N <sub>2</sub> /graphite	$\cong 140$		$\cong 250$
[125]	2500	tungsten/dry N <sub>2</sub> /mica	$230 \pm 30$		$330 \pm 150$
		/graphite	$140 \pm 90$		$220 \pm 120$
		/PTEE <sup>(c)</sup>	$2 \pm 2$	$21 \pm 2$ (18)	$5 \pm 5$
	2000	/stearic acid	$17 \pm 11$	$24 \pm 2$ (21)	$35 \pm 11$
		/Al <sub>2</sub> O <sub>3</sub>	$85 \pm 25$	$41 \pm 4$ (45)	$100 \pm 25$
	3000	/trifluorostearic acid	$5 \pm 4$	$23 \pm 2$ (20)	$20 \pm 8$
[55]	200	Ni/dry N <sub>2</sub> /Au	$\cong 3000$		4000–5000
[21]	270, 27.5	tungsten/air/graphite		15–40, 60–80	$\cong 24, \cong 8$
[95]	40	Si <sub>3</sub> N <sub>4</sub> /air/mica	0.6		
[126]	20–50	Si <sub>3</sub> N <sub>4</sub> /air/glass			300
[104]	100	Si <sub>3</sub> N <sub>4</sub> /air/mica			192 (184)
		gold/air/mica			51 (48)
		paraffin/air/mica			17 (16)
[62]	5000 <sup>(b)</sup>	iron oxide/air/silica		3 (163)	

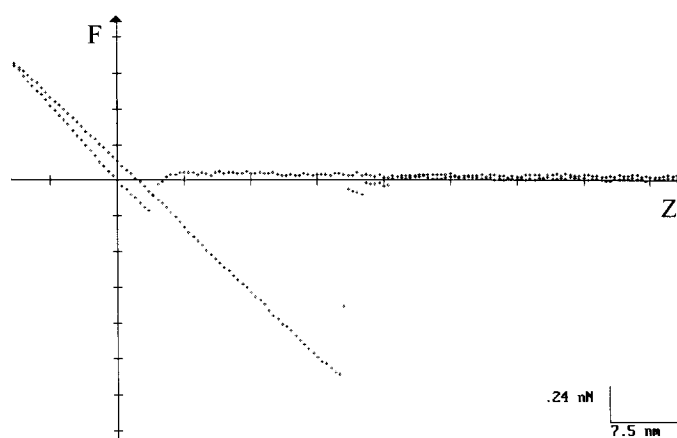


Fig. 24. Force–displacement curve on mica in water with a silicon nitride tip ( $R$  between 50 and 100 nm). Note the inverse path effect on the contact lines (adapted from [128]).

Table 4

Measured attractive force  $F_{\text{attr}}$ , adhesion force  $F_{\text{ad}}$ , and Hamaker constants  $A$  for experiments in water. Tips labeled with (a) are colloidal probes.

Ref.	Tip radius (nm)	Tip material/Medium/Sample	Measured (calculated)		
			$A$ ( $\times 10^{-19}$ J)	$F_{\text{attr}}$ (nN)	$F_{\text{ad}}$ (nN)
[91]	9000 <sup>(a)</sup>	TiO <sub>2</sub> /TiO <sub>2</sub>	0.6 $\pm$ 0.2		
[93]	3300 <sup>(a)</sup> , 100 <sup>(a)</sup>	Au/Au	2.5, 2.5		
[69]	160–260 <sup>(a)</sup>	Si <sub>3</sub> N <sub>4</sub> /Si <sub>3</sub> N <sub>4</sub>	0.32–2.2 (0.61)		
	130–280 <sup>(a)</sup>	Si <sub>3</sub> N <sub>4</sub> /mica	0.06–1.7 (0.34)		
[127]	30–50	Si <sub>3</sub> N <sub>4</sub> /stearyl amine			3–5
		/stearic acid			1–1.5
		/stearyl amide			$\cong$ 0.5
[48]	20	Si <sub>3</sub> N <sub>4</sub> /mica			0.8–1.4
		/lysozyme			0.2–0.4
[77]	220 <sup>(a)</sup>	Si <sub>3</sub> N <sub>4</sub> /Si <sub>3</sub> N <sub>4</sub>		1.3	
[101]	50	Si <sub>3</sub> N <sub>4</sub> /SiO <sub>2</sub>			44 (45)
		/Si <sub>3</sub> N <sub>4</sub>			50
[62]	5000 <sup>(a)</sup>	iron oxide/silica			( $\gamma$ in mJ/m <sup>2</sup> )
					0.3 (51)
[128]	$\cong$ 50	Si <sub>3</sub> N <sub>4</sub> /mica		0.24	1.4
		Si <sub>3</sub> N <sub>4</sub> /stearic acid		1–3	
		Si <sub>3</sub> N <sub>4</sub> /Au		0.4	10

Hartmann [123, 129] has calculated, according to Lifshitz theory, the Van der Waals force in different liquids (H<sub>2</sub>O, glycerol, CCl<sub>4</sub>, H<sub>2</sub>O<sub>2</sub>, formamide, and glycol) verifying that the force is reduced by a 10-fold factor (or 100 for the last three liquids). For CCl<sub>4</sub> the Van der Waals force is attractive at any distance. For formamide and glycerol it is repulsive, because their dielectric functions exhibit a rather strong polar contribution. For H<sub>2</sub>O, H<sub>2</sub>O<sub>2</sub>, and glycol there is a transition point from the attractive to the repulsive regime between 100 and 1000 nm. This effect is due to the retardation of the dispersion contribution that becomes smaller than the polar repulsive term. Unfortunately, the Van der Waals force becomes very small at such distances. Hence it turns out to be very difficult to detect such inversion points. In Fig. 25 a force–displacement curve acquired on mica in formamide is shown. Before the jump-to-contact and after the jump-off-contact, the force follows a  $D^{-2}$  law and is repulsive.

In water, another phenomenon can be observed, namely the hydrophobic meniscus force. If the sample surface is made hydrophobic by means of a layer of hydrocarbons and is dipped in water, after the tip has contacted the sample, hydrocarbon molecules form a meniscus that exerts a strong adhesion force. Thus the sample has to be withdrawn for a rather long distance before the tip detaches from the sample surface. Such a meniscus force results in withdrawal curves with a “slide-off-contact”, i.e., a gradual and continuous pull-off, instead of a jump-off-contact. The attractive hydrophobic force can be hardly seen in Fig. 26. The jump-to-contact is rather big, however, and this cannot be due to the Hamaker constant. The attractive hydrophobic force is treated in detail in Section 6.6. Similar curves can be obtained on gold or platinum [61], because metals become hydrophobic when exposed to air due

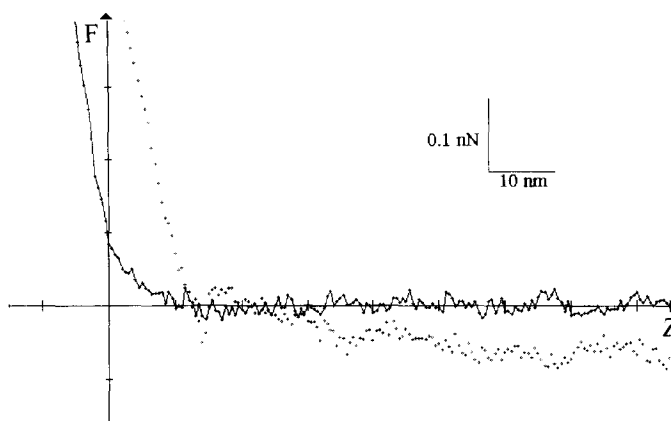


Fig. 25. Force–displacement curve in formamide with the repulsive van der Waals force just before the contact line and just past jump-off-contact (the continuous line is the approach curve). Note the splitting of the two zero lines of the approach and the retraction curve due to the viscosity of the medium (adapted from [128]).

to the adsorption of a thin layer of alkanes. Such “slide-off” curves become “jump-off” curves if gold is dipped into ethanol or cleaned with ethanol, which dissolves the alkane layer.

Experiments carried out in several liquids are summarized in Table 5. In addition to the works listed in Tables 4 and 5, those on double-layer force also contain measurements of the Van der Waals force of the system, since it is necessary to subtract the Van der Waals contribution in order to correctly estimate the double-layer force. Sometimes the Van der Waals force is computed on the basis of theoretical Hamaker constants in order to measure experimentally the double-layer force. Some authors compute the double-layer force and then experimentally measure the Van der Waals force.

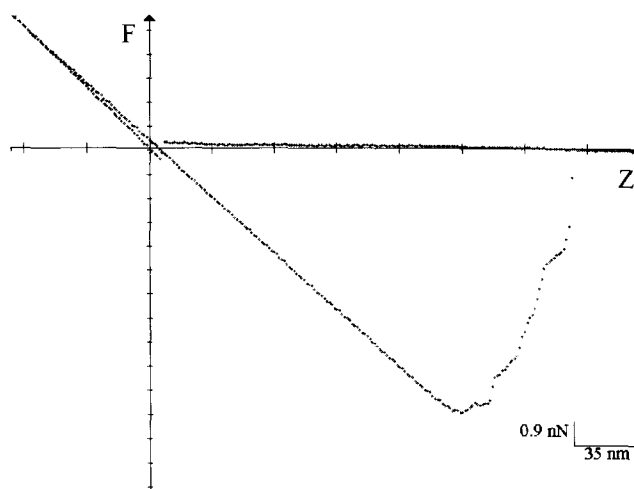


Fig. 26. Hydrophobic meniscus on gold in water. The pull-off is larger than on mica, and it is gradual and continuous over a distance of about 40 nm. Also the pull-off distance is larger than on mica (adapted from [128]).



Table 5

Measured attractive force  $F_{\text{attr}}$ , adhesion force  $F_{\text{ad}}$ , and surface energy  $\gamma$  for experiments in different liquids. Quantities labeled with (a) are unknown, tips labeled with (b) are colloidal particle. (c) PTFE = poly-tetra-fluoro-ethylene, PFMCH = per-fluoro-methyl-cyclo-hexane.

Ref.	Tip radius (nm)	Tip material/Medium/Sample	Measured (calculated)		
			$A$ ( $\times 10^{-19}$ J)	$F_{\text{attr}}$ (nN)	$F_{\text{ad}}$ (nN)
[61]	50	Si <sub>3</sub> N <sub>4</sub> /ethanol/mica, Au		0.14	0.2, 0.2–0.5
		/mica		0.06	0.2–1.5
		/platinum		26	55
		/H <sub>2</sub> O + ethanol/platinum		20	29
		/ethanol/platinum		0	0
		/formamide/mica		– 0.12	– 0.12
		(a) WC/ethanol/Au		0.3	0.5
		/mica			0.2–0.5
		/mica			0.5–2
		/Au		6	10
[78]	50	/formamide/Au		7	15
		Si <sub>3</sub> N <sub>4</sub> /ethanol/mica	1.16		
		/bromo-naphtalene/mica	– 0.07		
[101]	50	/methyl-naphtalene/mica	– 2.36		
		Si <sub>3</sub> N <sub>4</sub> /acetone/SiO <sub>2</sub> , Si <sub>3</sub> N <sub>4</sub>			14 (15), 32
[131]	(a), (b)	/ethanol/SiO <sub>2</sub> , Si <sub>3</sub> N <sub>4</sub>			12 (13), 20
		Au/cyclohexane/PTFE <sup>(c)</sup>	– 0.55 (– 0.52)		
		/p-Xylene/	– 0.5 (– 0.48)		
		/bromobenzene/	– 0.55 (– 0.76)		
		/perfluorohexane/	0.07 (0.125)		
		/PFMCH <sup>(c)</sup> /	0.05		
		/formamide	– 0.4		
[128]	$\cong$ 50	Si <sub>3</sub> N <sub>4</sub> /formamide/mica		– 0.1	
		/ethanol/		0.1	
[130]	50	Si <sub>3</sub> N <sub>4</sub> /diodomethane/Si <sub>3</sub> N <sub>4</sub> , SiO <sub>2</sub>	(0.1), (– 0.08)	0.1, $\cong$ – 0.2	0.6, $\cong$ – 0.2
		/bromonaphtalene/SiO <sub>2</sub> , Si <sub>3</sub> N <sub>4</sub>	(– 0.02), (0.3)	$\cong$ – 0.1	$\cong$ – 0.1

The work of Weisenhorn et al. [61] has been of outstanding importance in the development of force–distance curves acquisition. Measurements have been done in water, ethanol and formamide for three kinds of systems: non-metal/non-metal, non-metal/metal and metal/metal. In water, the forces for non-metal/metal or metal/metal systems (Si<sub>3</sub>N<sub>4</sub> on platinum or gold and WC on Au) are very much bigger than for non-metal/non-metal systems (Si<sub>3</sub>N<sub>4</sub> on mica). The authors demonstrate that this is due to the hydrophobic meniscus by dipping the metals in ethanol or ethanol–water mixtures. The force progressively decreases due to the dissolving of the alkane layer. Forces in ethanol are of the same order for all the three systems. In formamide, the force between Si<sub>3</sub>N<sub>4</sub> and mica is repulsive, while it is strongly attractive for WC and platinum. This is once again due to the hydrophobic meniscus, which is smaller than in water.

Table 6

Measured and calculated adhesion force for different polymers in perfluorodecalin. All data are taken from the Ref. [132]. The polystyrene tip is a colloidal particle. (a) FEP = poly-tetra-fluoro-ethylene-co-hexa-fluoro-propylene; PVDF = poly-vinylidene-fluoride; PAN = poly-acrylonitrile; PAA = poly-acrylic acid; PMMA = poly-(methyl-methacrylate).

Tip radius (nm)	Tip material/Sample	Measured (calculated) $F_{ad}$ (nN)
51	SiO <sub>2</sub> /FEP <sup>(a)</sup>	$0.18 \pm 0.08$ (0.39)
	/PVDF <sup>(a)</sup>	$0.62 \pm 0.2$ (1.23)
	/polypropylene	$2.07 \pm 0.15$ (1.98)
	/polystyrene	$2.98 \pm 0.16$ (2.76)
20	/PAN <sup>(a)</sup>	$1.32 \pm 0.15$ (0.26)
	/PMMA <sup>(a)</sup>	$1.84 \pm 0.16$ (0.73)
	/PAA <sup>(a)</sup>	$2.13 \pm 0.14$ (0.84)
	/polystyrene	$0.66 \pm 0.1$ (1.08)
3750	polystyrene/PAN <sup>(a)</sup>	$54.4 \pm 20.6$ (55.5)
	/PMMA <sup>(a)</sup>	$113 \pm 4$ (203)
	/PAA <sup>(a)</sup>	$75.4 \pm 4.5$ (231)
	/polystyrene	$135 \pm 7.5$ (314)

In the work of Hutter and Bechhoefer [78] different liquids are chosen in order to minimize the Hamaker constant and hence the Van der Waals attraction. The Van der Waals force is attractive in ethanol and repulsive in methyl-naphthalene, but it is nearly zero in bromo-naphthalene. Therefore, this liquid could be used in contact mode to enhance lateral resolution (see also Ref. [130]).

Finally, the work of Milling et al. [131] is designed to measure repulsive Van der Waals forces in different liquids. Nevertheless, some liquids with attractive Van der Waals forces are studied.

Feldman et al. [132] have shown the possibility of a kind of force spectroscopy on polymers. The best liquid for distinguishing samples is perfluorodecalin (PFD), since it shows the best signal-to-noise ratio. Measurements are all performed in PFD (see Table 6) using a glass or a polystyrene colloidal tip. The “exact” adhesion force is calculated on the basis of the JKR theory. Although a sensible trend for the adhesion force is found in the case of the first four polymers listed in Table 6, considerable discrepancies between the calculated and measured values occur for all the other measurements. These discrepancies are attributed to the microasperities of the polystyrene tip and to the presence of other interactions (e.g., hydrogen bonding) not accounted for in the Lifshitz theory.

A very interesting experiment is that of Mulvaney et al. [133], in which the authors have measured the forces between an oil (decane) droplet and a glass colloidal particle. The contact lines of the curves acquired with such a system depend on the scan rate. At very low scan rates ( $<0.5$  Hz), the lines show a large number of high frequency oscillations and in the region between 0.5 and 2 Hz the lines show numerous, regularly spaced waves. At higher scan rates the oscillations are no longer evident. The Van der Waals force between glass and decane in water has been measured at low scan rates and the Hamaker constant is  $6 \times 10^{-21} \pm 3 \times 10^{-21}$  J, in good agreement with the theoretical value of  $3 \times 10^{-21} - 7 \times 10^{-21}$  J.

The experiments dealing with the determination of the Van der Waals force have demonstrated that the AFM is a useful tool in the study of this fundamental force. Measurements can be performed in any kind of liquid and between any kind of surfaces with high resolution. The exact determination of Hamaker constants is not a simple issue, but for several materials this has been achieved in good agreement with theoretical predictions. In other cases, the presence of other forces, the difficulty of eliminating contaminants, and the lack of knowledge of the tip shape make the measurement more problematic. The capability of distinguishing the Hamaker constants of very similar materials also has been demonstrated. The recent measurements of the Van der Waals force at a liquid/liquid interface are of outstanding importance.

### 6.2.3. Dependence on tip shape

Hartmann [111] has calculated the effect of tip shape assuming a Van der Waals force in the form:

$$F = g \frac{A}{D^n}, \quad (6.36)$$

in which  $g$  is a geometric factor that he has calculated by means of Lifshitz theory and  $A$  is the Hamaker constant, which is shape independent.

Hartmann has pointed out that, like Hamaker constants, geometric factors also depend on retardation. The retarded Van der Waals force decays on smaller distances, and hence involves a smaller effective volume than the non-retarded Van der Waals force. The non-retarded factor  $g_n$ , the retarded one  $g_r$ , and the power law index  $n$  of the non-retarded Van der Waals force are indicated in Table 7. The power law index of the retarded force is  $n + 1$ , and there is a transition from  $n$  to  $n + 1$  for  $D \geq 5\lambda_{132}$ . In the present context the most important result is that the Van der Waals force decays as  $D^{-2}$  for the most realistic probe geometry, i.e., the paraboloidal apex.

Hao et al. [21] have calculated a correction factor for the flat-sphere system in order to take into account the elongated nature of the tip. Considering the tip as a cone of semiaperture  $\theta/2$  and length  $L$ , ending with a sphere of radius  $R$ , for  $D \ll L$  they obtain:

$$F(D) = -\frac{AR}{6D^2} \left( 1 + \frac{D \operatorname{tg}^2(\theta/2)}{R} \right). \quad (6.37)$$

The correction factor is negligible for  $D \ll R$ .

The distribution of microasperities on the tip and/or on the sample has a dramatic effect on the calculated Van der Waals forces. Generally, the effect of microasperities is to greatly reduce the force,

Table 7

Retarded and non-retarded geometric factors for three geometries as calculated by Hartmann [111].  $l_{xy}$  and  $l_z$  are the semiaxes of the paraboloid,  $R$  is the radius of the cylinder and  $\theta$  the semiaperture of the cone.

	$g_n$	$g_r$	$n$
Cone	$(\operatorname{tg}^2\theta)/6$	$(\pi/3)\operatorname{tg}^2\theta$	1
Paraboloid	$l_{xy}^2/12l_z$	$\pi l_{xy}^2/3l_z$	2
Cylinder	$R^2/6$	$\pi R^2$	3

because the effective contact radius is smaller than the overall radius. Cohen et al. [43] have calculated that a distribution of microtips with mean radius of 20 Å and 4 Å rms deviation would reduce the adhesion force by more than 90%.

Bordag et al. [134] have calculated the corrections to the Van der Waals force corresponding to two types of deviations of the tip from the paraboloidal shape. The deviations of the first type are small protrusions on the tip surface that, in this case, can be described by the formula:

$$z_1(\rho) = \frac{\rho^2}{2R} + \xi f(x, y), \quad (6.38)$$

in which  $\rho$  is the cylindrical coordinate,  $R$  the tip radius, and  $\xi f(x, y)$  is the stationary stochastic function with dispersion  $\xi$  and mean value  $\langle \xi f(x, y) \rangle = 0$ . The second type of deviations corresponds to a possible deformation of the tip surface described by

$$z_2 = \frac{[\rho + \xi g(\varphi)]^2}{2R}, \quad (6.39)$$

in which  $\rho$  and  $\varphi$  are cylindrical coordinates and  $g(\varphi)$  is another stationary stochastic function with dispersion  $\xi$  and zero mean value.

In the first case the correction to the Van der Waals force is

$$F = -\frac{AR}{6D^2} \left[ 1 + 3\left(\frac{\xi}{D}\right)^2 + 15\left(\frac{\xi}{D}\right)^4 \right]. \quad (6.40)$$

In the second case the correction is

$$F = -\frac{AR}{6D^2} \left[ 1 + \left(\frac{\xi}{\sqrt{RD}}\right)^2 - \frac{3}{4} \left(\frac{\xi}{\sqrt{RD}}\right)^4 \right]. \quad (6.41)$$

For  $\xi/D \cong 0.1$  the contribution of the second order correction is about 3%, while the contribution of the fourth order term is negligible. For  $\xi/\sqrt{RD} \cong 0.1$  the correction to  $F_0$  according to Eq. (6.41) is about 0.1%.

Toikka et al. [62] have verified the effect of microasperities on the tip. They measured the Van der Waals force for the system iron oxide/air or water/silica with a colloidal probe of radius  $\cong 5\mu\text{m}$ . The measured value of the Van der Waals adhesion energy is very much smaller than the calculated value (see Table 2). If the calculation is made with the characteristic dimension of the asperities (between 10 and 40 nm) instead of the overall radius of the colloidal probe, experimental and calculated values are in agreement. They argue that asperities on the tip decrease the adhesion force not only because the effective radius is smaller than the overall radius, but also because a thin layer of water solution stays between the tip and the sample surface, thus exerting a repulsive double-layer force.

Cohen [39] has considered the effect of a layer of contaminants of height  $X$  adsorbed on a tip of radius  $R$  at a distance  $D$  from the sample surface. The interaction between the spherical tip and the flat sample is affected by the layer of contaminants as follows:

$$W = \frac{A}{6} \left[ \frac{2RD + D^2}{2(D + X)^2} - \frac{2(D + R)}{(D + X)} + \frac{3D + 2R}{2D} + \ln D + \ln(D + X) \right]. \quad (6.42)$$

For  $R \gg D$ , the force is given by

$$F = -\frac{A}{6} \left[ \frac{R}{D^2} - \frac{RD + 3RX}{(D + X)^3} \right]. \quad (6.43)$$

For  $D = 5 \text{ \AA}$  and  $X = 7 \text{ \AA}$  the contaminant layer decreases the attractive force by about 50%.

### 6.3. Double-layer force

#### 6.3.1. Theory

Usually, all surfaces of high dielectric constant in water or other liquids are charged. The charging may occur via either dissociation of surface groups (e.g.,  $\text{COOH} \rightarrow \text{COO}^- + \text{H}^+$ ) or adsorption of ions from the solution. Whatever be the charging mechanism, the final surface charge is balanced by a charged region of counterions of equal and opposite charge. Some of these counterions are bound to the surface in the so-called Stern layer, while others form an atmosphere of ions near the surface known as electric double-layer.

Following [15], consider two equally charged surfaces at a distance  $D$  and a liquid solution in between. The coordinate system is illustrated in panel (a) of Fig. 27.

If ions are present in solution, the chemical potential of each species of ion can be written as

$$\mu = ze\psi(x) + k_B T \log \rho(x), \quad (6.44)$$

in which  $\psi(x)$  is the electrostatic potential and  $\rho(x)$  is the number density of the counterions of valence  $z$  at each point  $x$  between the surfaces. Let  $\psi(0) = \psi_0 = 0$ ,  $\rho(0) = \rho_0$ , and  $(d\psi/dx)_0 = 0$  ( $x = 0$  is the midplane between the two surfaces). Then the Boltzmann distribution of counterions is given by

$$\rho(x) = \rho_0 \exp(-ze\psi(x)/k_B T). \quad (6.45)$$

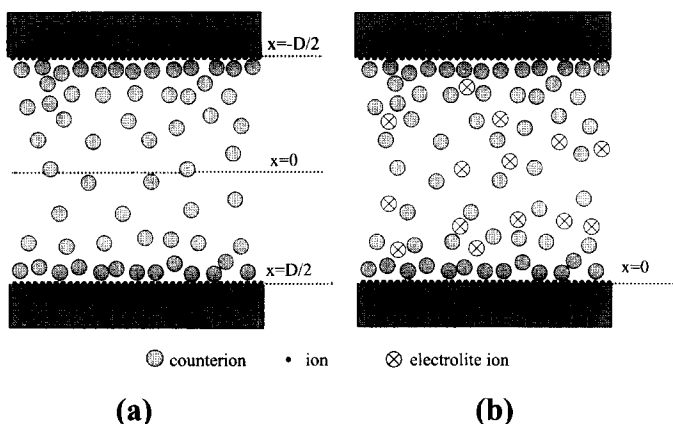


Fig. 27. Panel (a): the coordinate system employed in the derivation of the double-layer force without electrolyte ions (Eqs. (6.44)–(6.59)).  $D$  is the tip–sample distance and the tip and sample surfaces are at  $x = \pm D/2$ .  $x = 0$  is the midplane. Panel (b): the coordinate system employed in the derivation of the double-layer force with electrolyte ions (Eqs. (6.60)–(6.75)).  $D$  is the tip–sample distance.  $x = 0$  is the sample surface. The distribution of the counterions and of the electrolyte ions in the liquid gap is also sketched.

Eq. (6.45), together with the Poisson equation, yields

$$\frac{d^2\psi(x)}{dx^2} = -\frac{ze\rho(x)}{\varepsilon\varepsilon_0} = -\frac{ze\rho_0}{\varepsilon\varepsilon_0} \exp\left(-\frac{ze\psi(x)}{k_B T}\right). \quad (6.46)$$

This is the Poisson–Boltzmann equation. Solving Eq. (6.46), one obtains the potential  $\psi(x)$ , the electric field  $E(x) = d\psi(x)/dx$ , and the counterion density  $\rho(x)$ .

The two boundary conditions needed to solve Eq. (6.46) are

$$E_0 = \left. \frac{d\psi(x)}{dx} \right|_0 = 0, \quad (6.47a)$$

$$\sigma = -\int_0^{D/2} ze\rho(x) dx = \varepsilon\varepsilon_0 \int_0^{D/2} \frac{d^2\psi(x)}{dx^2} dx = \varepsilon\varepsilon_0 \left. \frac{d\psi(x)}{dx} \right|_{D/2} = \varepsilon\varepsilon_0 E_s. \quad (6.47b)$$

The latter condition means that the total charge in the liquid gap is equal and opposite to the surface charge,  $\sigma$ .

Next consider the counterions distribution. Differentiating Eq. (6.45) and inserting Eq. (6.46), yields

$$\frac{d\rho(x)}{dx} = -\frac{ze\rho_0}{k_B T} \exp\left(-\frac{ze\psi(x)}{k_B T}\right) \frac{d\psi(x)}{dx} = \frac{\varepsilon\varepsilon_0}{k_B T} \frac{d\psi(x)}{dx} \frac{d^2\psi(x)}{dx^2} = \frac{\varepsilon\varepsilon_0}{2k_B T} \frac{d}{dx} \left( \frac{d\psi(x)}{dx} \right)^2 \quad (6.48)$$

and hence

$$\rho(x) - \rho_0 = \int_0^x d\rho = \frac{\varepsilon\varepsilon_0}{2k_B T} \int_0^x d \left( \frac{d\psi(x)}{dx} \right)^2 = \frac{\varepsilon\varepsilon_0}{2k_B T} \left( \frac{d\psi(x)}{dx} \right)^2 \Big|_x. \quad (6.49)$$

Using Eq. (6.47b), the value  $\rho_s$  of  $\rho(x)$  at the surface ( $x = \pm D/2$ ) can be determined:

$$\rho_s = \rho_0 + \frac{\sigma^2}{2\varepsilon\varepsilon_0 k_B T}. \quad (6.50)$$

Eq. (6.50) shows that the concentration of ions at the surface depends only on  $\sigma$  and  $\rho_0$ . Most of the counterions balancing the surface charge are placed in the first few angstroms near the surface. However, for lower densities  $\sigma$ , since  $\rho_s \propto \sigma^2$ , the counterions layer extends well beyond the surface.

The Poisson–Boltzmann equation can be satisfied by

$$\psi(x) = \frac{k_B T}{ze} \log \cos^2 \Gamma x, \quad \Gamma = \sqrt{\frac{(ze)^2 \rho_0}{2\varepsilon\varepsilon_0 k_B T}}. \quad (6.51)$$

$\Gamma$  can be determined from Eq. (6.51) if  $E_s$  is known:

$$E(x) = \frac{d\psi(x)}{dx} = -\frac{2k_B T \Gamma}{ze} \operatorname{tg} \Gamma x \Rightarrow E_s = \frac{\sigma}{\varepsilon\varepsilon_0} = -\frac{2k_B T \Gamma}{ze} \operatorname{tg} \frac{\Gamma D}{2}. \quad (6.52)$$

The counterion distribution

$$\rho(x) = \rho_0 \exp\left(-\frac{ze\psi(x)}{k_B T}\right) = \frac{\rho_0}{\cos^2 \Gamma x} \quad (6.53)$$

is known once  $\Gamma$  is determined in terms of  $\sigma$  and  $D$ .

Since  $\partial P/\partial x|_{x,T} = \rho \partial \mu/\partial x|_{x,T}$ , the variation in pressure resulting from bringing the two surfaces from an infinite distance to a distance  $x = D/2$  is

$$P(x) = - \int_{\xi=D/2}^{\xi=\infty} \left[ ze \rho \frac{d\psi(x)}{dx} \right]_x d\xi + k_B T d\rho(x), \quad (6.54)$$

in which  $x$  is the fixed point in which values are computed and  $\xi$  is the variable separation between surfaces. Substituting the Poisson equation into the previous equation we obtain

$$P^D(x) - P^\infty(x) = -\frac{\varepsilon \varepsilon_0}{2} \left( \frac{d\psi(x)}{dx} \Big|_{x(D)} \right)^2 + k_B T \rho^D(x) + \frac{\varepsilon \varepsilon_0}{2} \left( \frac{d\psi(x)}{dx} \Big|_{x(\infty)} \right)^2 - k_B T \rho^\infty(x), \quad (6.55)$$

in which  $P^D(x)$ ,  $P^\infty(x)$ ,  $\rho^D(x)$ , and  $\rho^\infty(x)$  are the pressures and the distributions when the two surfaces are at a distance  $D$  and at an infinite distance. In the case that no electrolyte ions are present in solution,  $\rho^\infty(0) = P^\infty(x) = 0$ . Thus  $P^D(x)$  is the sum of two contributions. The first one, always attractive, is the contribution of the electrostatic field and the second one is the osmotic or entropic contribution and it is always repulsive.

From Eq. (6.55) one can see that the origin of the repulsive force between charged surfaces in solution is not electrostatic (the electrostatic term is always attractive), but entropic. What keeps the double-layer of counterions expanded is the osmotic pressure that compels counterions away from the surfaces and far from each other, in order to increase the configurational entropy. By bringing the two surfaces in contact, the counterions in solution are forced to come in contact with the surfaces, against their equilibrium configuration.

Inserting Eqs. (6.49) and (6.50) into Eq. (6.55) it is possible to derive an expression for the pressure at the surface ( $x = \pm D/2$ ):

$$P_s^D = k_B T \rho^D(0) = k_B T [\rho_s^D - \rho_s^\infty]. \quad (6.56)$$

This important equation, known as the contact value theorem, is always valid as long as there is no interaction between counterions and surfaces.

Using Eqs. (6.46) and (6.51) to evaluate  $\rho^D(0)$  in Eq. (6.56) leads to an expression for the pressure in terms of  $\Gamma$ :

$$P = k_B T \rho_0 = 2\varepsilon \varepsilon_0 \left( \frac{k_B T}{ze} \right)^2 \Gamma^2. \quad (6.57)$$

In order to evaluate Eq. (6.57) for two limiting cases  $\sigma \rightarrow \infty$  and  $D \rightarrow \infty$  one needs to know  $\Gamma$ . For  $\sigma \rightarrow \infty$ , Eq. (6.52) yields  $\Gamma D/2 \rightarrow \pi/2$ . For  $D \rightarrow 0$ ,  $\text{tg } \Gamma D/2 \cong \Gamma D/2$  and Eq. (6.52) yields  $\Gamma^2 \rightarrow -\sigma ze/\varepsilon \varepsilon_0 k_B T D$ . In the first case the Langmuir equation is obtained:

$$\lim_{\Gamma D/2 \rightarrow \pi/2} P(D) = \frac{2\varepsilon \varepsilon_0 (\pi k_B T / ze)^2}{D^2}. \quad (6.58)$$

In the second case:

$$\lim_{\Gamma^2 \rightarrow -\sigma ze/\varepsilon \varepsilon_0 k_B T D} P(D) = -\frac{2\sigma k_B T}{ze D}. \quad (6.59)$$

The simple form for the Poisson equation, Eq. (6.46), is valid only if there are no electrolyte ions in solution. If there are electrolyte ions in solution, Eq. (6.45) must be written for each kind of ion and the sum used for  $\rho(x)$  in Eq. (6.46). In this case the total concentration of ions at the surface is

$$\sum_i \rho_{0i} = \sum_i \rho_{\infty i} + \frac{\sigma^2}{2\varepsilon\varepsilon_0 k_B T} \quad (6.60)$$

in which  $z_i$  is the valence,  $\rho_{0i}$  the distribution at the surface, and  $\rho_{\infty i}$  is the distribution at an infinite distance for the  $i$ th ion. In the following, the  $x$  coordinate origin is changed, since one of the surfaces is now at infinite distance. Hence  $x = 0$  now indicates the surface position and no longer the midplane. The new system of coordinate is illustrated in panel (b) of Fig. 27.

The ions at the surface are mainly the counterions. Their excess concentration over the bulk concentration depends only on  $\sigma$ . From Eq. (6.60), it is possible [15] to determine the relation between  $\sigma$  and  $\psi_0$  in the case of a mixed KCl+CaCl<sub>2</sub> electrolyte:

$$\begin{aligned} \sigma^2 &= 2\varepsilon\varepsilon_0 k_B T \left( \sum_i \rho_{0i} - \sum_i \rho_{\infty i} \right) \\ &= 2\varepsilon\varepsilon_0 k_B T ([K^+]_{\infty} \exp(-e\psi_0/k_B T) + [Ca^{2+}]_{\infty} \exp(-2e\psi_0/k_B T) + [Cl^-]_{\infty} \exp(e\psi_0/k_B T)) \\ &\quad - 2\varepsilon\varepsilon_0 k_B T ([K^+]_{\infty} + [Ca^{2+}]_{\infty} + [Cl^-]_{\infty}). \end{aligned} \quad (6.61)$$

Since  $[Cl^-] = [K^+] + 2[Ca^{2+}]$ , Eq. (6.61) becomes

$$\begin{aligned} \sigma &= \sqrt{2\varepsilon\varepsilon_0 k_B T} [K^+]_{\infty} \{ \exp(-e\psi_0/k_B T) + \exp(e\psi_0/k_B T) - 2 \} \\ &\quad + \sqrt{2\varepsilon\varepsilon_0 k_B T} [Ca^{2+}]_{\infty} \{ \exp(-2e\psi_0/k_B T) + 2 \exp(e\psi_0/k_B T) - 3 \} \\ &= \sqrt{8\varepsilon\varepsilon_0 k_B T} \sinh(e\psi_0/k_B T) \sqrt{[K^+]_{\infty} + [Ca^{2+}]_{\infty} (2 + \exp(-e\psi_0/k_B T))}. \end{aligned} \quad (6.62)$$

Eq. (6.62) is the Grahame equation. This equation shows that, at constant surface charge density, the surface potential decreases with increasing electrolyte concentration. In the presence of divalent cations the surface potential decreases more than in the case of monovalent ions.

In most cases, when there are electrolyte ions in solution, neither  $\sigma$  nor  $\psi_0$  stay constant, because of reactions at the surface. Such reactions at the surface can be described by means of a reaction constant that can be inserted in Grahame equation. Parsegian and Gingell [135] have studied the consequences of these surface reactions. Their results can be summarized as follows:

- (1) The interaction potential is always intermediate between two limits.
- (2) The upper limit corresponds to the case of constant charge.
- (3) The lower limit corresponds to the case of constant electric potential.

A particularly interesting limit of the Grahame equation is the limit at small potentials. Below 25 mV, Eq. (6.62) becomes [15]:

$$\sigma = \varepsilon\varepsilon_0 K_D \psi_0, \quad (6.63)$$

in which

$$K_D = \sqrt{\sum_i \frac{\rho_{\infty i} e^2 z_i^2}{\varepsilon\varepsilon_0 k_B T}}. \quad (6.64)$$



Eq. (6.63) is analogous to the equation for a capacitor with charge density  $\pm\sigma$ , potential  $\psi_0$ , and whose plates are separated by a distance  $1/K_D$ . The characteristic length  $1/K_D$  is the Debye length. The Debye length depends only on solution properties and not on surface properties. At 25 °C the Debye length of aqueous solutions is:

$$\frac{1}{K_D} = \begin{cases} 0.304/\sqrt{[X]} \text{ nm} & \text{for 1:1 electrolytes,} \\ 0.176/\sqrt{[X]} \text{ nm} & \text{for 1:2 or 2:1 electrolytes,} \\ 0.152/\sqrt{[X]} \text{ nm} & \text{for 2:2 electrolytes,} \end{cases} \quad (6.65)$$

in which  $[X]$  is the electrolyte concentration in moles.

In the general case, i.e. for any potential, the electrostatic potential is given by

$$\psi(x) = \frac{2k_B T}{e} \log \frac{1 + \gamma e^{-K_D x}}{1 - \gamma e^{-K_D x}} \cong \frac{4k_B T}{e} \gamma e^{-K_D x}, \quad (6.66)$$

where

$$\gamma = \tanh \frac{e\psi_0}{4k_B T}. \quad (6.67)$$

This is known as the Gouy–Chapman theory. At strong electric potentials  $\gamma \rightarrow 1$ , whereas, for small potentials, Eq. (6.66) becomes the so-called Debye–Hückel equation:

$$\psi(x) \cong \psi_0 e^{-K_D x}. \quad (6.68)$$

For an 1:1 electrolyte the interaction pressure and the interaction energy are given by [15]:

$$P = 64k_B T \rho_\infty \gamma^2 e^{-K_D D} \quad \text{and} \quad W = 64k_B T \rho_\infty \frac{\gamma^2}{K_D} e^{-K_D D}. \quad (6.69)$$

For small potentials Eq. (6.69) becomes

$$P = 2 \frac{\sigma^2}{\varepsilon \varepsilon_0} e^{-K_D D} \quad \text{and} \quad W = 2 \frac{\sigma^2}{K_D \varepsilon \varepsilon_0} e^{-K_D D}. \quad (6.70)$$

Since  $K_D$  is proportional to the electrolyte concentration  $\rho$ , the force and the potential are inversely proportional to the salt concentration.

Generally in an AFM the tip and sample have different surface charge densities. Parsegian and Gingell [135] have described the double-layer force per unit area between two semiinfinite surfaces with surface charge density  $\sigma_1$  and  $\sigma_2$  at a distance  $D$  in the limit of small potentials ( $< 25$  mV):

$$f_{dl} = \frac{2}{\varepsilon \varepsilon_0} \frac{\sigma_1^2 + \sigma_2^2 + \sigma_1 \sigma_2 (e^{K_D D} + e^{-K_D D})}{(e^{K_D D} - e^{-K_D D})^2}. \quad (6.71)$$

If  $K_D D \gg 1$ , Eq. (6.71) becomes

$$f_{dl} = \frac{2}{\varepsilon \varepsilon_0} [(\sigma_1^2 + \sigma_2^2) e^{-2K_D D} + \sigma_1 \sigma_2 e^{-K_D D}]. \quad (6.72)$$

Butt [136] has computed the double-layer force between a flat sample and a spherical tip with radius  $R$  by integrating Eq. (6.72) on circular sections  $2\pi r dr$  (Derjaguin approximation):

$$F_{dl} = \frac{\pi}{\varepsilon\varepsilon_0 K_D^2} [(\sigma_T^2 + \sigma_S^2)(e^{-2K_D R} + 2K_D R - 1)e^{-2K_D D} + 4\sigma_T\sigma_S(K_D R + e^{-K_D R} - 1)e^{-K_D D}], \quad (6.73)$$

in which  $\sigma_T$  and  $\sigma_S$  are the tip and sample surface charge densities. If  $K_D R \gg 1$ , Eq. (6.73) becomes

$$F_{dl} = \frac{2\pi R}{\varepsilon\varepsilon_0 K_D} [(\sigma_T^2 + \sigma_S^2)e^{-2K_D D} + 2\sigma_T\sigma_S e^{-K_D D}]. \quad (6.74)$$

If  $\sigma_T \ll \sigma_S$ , i.e., if the tip is neutral, the first term in Eq. (6.74) dominates and the double-layer force decays with  $1/2K_D$ , whereas, if  $\sigma_T \cong \sigma_S$ , the second term dominates and the double-layer force decays with  $1/K_D$ . However, for  $D > (1/K_D) \ln((\sigma_T^2 + \sigma_S^2)/2\sigma_T\sigma_S)$ , the second term in Eq. (6.74) is always dominant. It should also be noted that the double-layer force between two bodies with different surface charge may be attractive, whereas the double-layer force between two surfaces with the same surface charge density is always repulsive.

A similar, approximate expression can be derived in the case of constant potential [137]:

$$F_{dl} = 4\pi\varepsilon\varepsilon_0 R K_D \psi_T \psi_S e^{-K_D D}. \quad (6.75)$$

Also in this case, the double-layer force may be attractive or repulsive depending on the sign of the potentials product.

The behavior of oxide-like materials, such as Si–OH, is particularly important for the AFM. For such materials, the potential is governed by the amphoteric character of surface sites, i.e., the groups can act as proton donors (Si–OH  $\rightarrow$  SiO $^-$ ) or proton acceptors (Si–OH  $\rightarrow$  SiOH $_2^+$ ). Therefore, both  $pK^+$  for donors and  $pK^-$  for acceptors must be defined and the point of zero charge (PZC) is equal to the quantity  $(pK^+ + pK^-)/2$ . The surface potential is positive (negative) when the pH of the solution is smaller (larger) than the PZC. If both tip and sample are oxide-like materials, two isoelectric points (IEP) can be defined when the potential becomes zero for either the tip or the sample. Hence the double-layer force is repulsive for pH values smaller than the smallest IEP or larger than the largest IEP, and is attractive in between.

### 6.3.2. Experimental results

In order to introduce the experimental results on double-layer force, a specific AFM measurement of these forces is presented first. In this experiment, the AFM tip is made of silicon nitride and the substrate is mica. The electrolyte is KCl and the Debye length is varied by varying the KCl concentration at constant pH. The force–displacement curves at different concentrations are depicted in Fig. 28.

Increasing the concentration up to 0.1 M, three changes occur:

1. the repulsive force becomes smaller and smaller;
2. its decay length diminishes;
3. the Van der Waals force appears.

At 0.1 M KCl concentration, the force–displacement curve resembles the one in deionized water (Fig. 23). Since deionized water is a weak electrolyte ( $10^{-7}$  M in both H $_3$ O $^+$  and OH $^-$ ), it should show

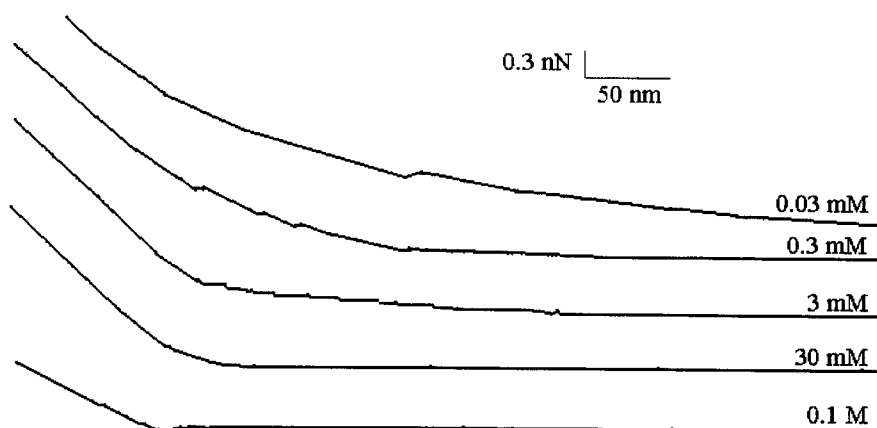


Fig. 28. Force–displacement approach curves associated with double-layer force between a silicon nitride tip and mica in KCl solutions of different concentration. The tip radius is estimated between 50 and 100 nm. On the x-axis, the sample position in nm. On the y-axis, the force in nN (adapted from [128]).

the highest double-layer force. However, this long range force is acting on distances well beyond the range of the AFM tip motion. Therefore, it is not possible to detect the actual zero line and to define the zero force reference.

These results were first obtained by Butt [138] with  $\text{Si}_3\text{N}_4$  tips on mica, diamond tips on mica and glass tips on glass, by Weisenhorn et al. [61], by Atkins and Pashley [92] with a ZnS colloidal sphere on mica, and by Li et al. [94] with polystyrene spheres on polystyrene spheres or mica. Only a qualitative comparison can be done, since the tip radius of curvature is unknown. In all papers, in order to deduce the Debye length, the force law has been fitted with the second term of Eq. (6.74), since one supposes  $K_D D \gg 1$  and  $K_D R \gg 1$ . A good agreement has been found in all the cases. Weisenhorn et al. have obtained a remarkable agreement with Eq. (6.65) (Debye length =  $0.308/\sqrt{[\text{KCl}]}$  nm). The fact that only the second term of Eq. (6.74) is needed to fit the experimental data implies that the surface charge density of diamond and silicon nitride are comparable to that of mica. This evidence contradicts what one would expect for these two materials. In the case of  $\text{Si}_3\text{N}_4$  the origin of the surface charge could be explained by the presence of small amounts of oxide [138].

Similar series of curves can be obtained in other electrolyte solutions, e.g.,  $\text{MgCl}_2$  [138],  $\text{CaCl}_2$  [139] or  $\text{Ca}(\text{NO}_3)_2$  [140]. Using a 2:1 salt, the Debye length decreases, and the double-layer force vanishes at a 30 mM concentration. In order to measure the dependence of double-layer force on tip and/or sample surface charge, a sample whose surface charge density depends on pH has to be employed.

Fig. 29 shows force–displacement curves on stearic acid in different pH solutions. Stearic acid pK is nearly 8. This means that, at pH less than 8, the sample surface is uncharged, there is no double-layer force and there is a jump-to-contact due to the Van der Waals force (1–3 nN). At pH 8.1 the surface begins to charge and the double-layer force counterbalances the Van der Waals force, so that there is no jump-to-contact. Finally, at a pH greater than 8, the double-layer force appears.

Ishino et al. [127] have characterized the charging of several functional groups as a function of pH. A behavior similar to that shown in Fig. 29 was observed for stearic acid and steryl amide (in this case, repulsive forces are observed for  $\text{pH} > 10.5$ ). Butt [138] has shown the transition from double-layer force to Van der Waals force for an alumina tip on mica. Mica is negatively charged at any pH value,

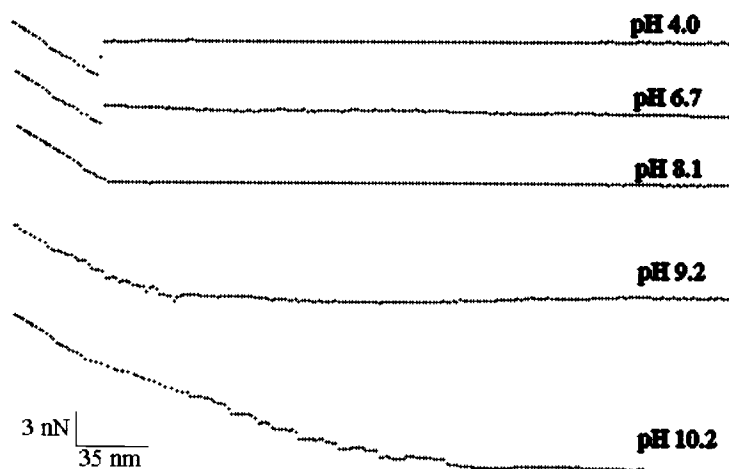


Fig. 29. Dependence of double-layer force on the surface charge of a stearic acid sample. Curves are acquired at different pH with a silicon nitride tip (tip radius between 50 and 100 nm). On the  $x$ -axis, the sample position in nm. On the  $y$ -axis, the force in nN.

whereas alumina is positively (negatively) charged below (above) pH 8.1. The forces are repulsive above pH 8.1 and attractive below pH 8.1.

Several studies have been dedicated to the double-layer force on oxide-like materials.

Raiteri et al. [141] have studied the pH dependence of forces for  $\text{Si}_3\text{N}_4$  tips on  $\text{Si}_3\text{N}_4$ ,  $\text{Al}_2\text{O}_3$ , and mica, demonstrating the capability of AFM in determining the PZC of such materials. For  $\text{Si}_3\text{N}_4$  on  $\text{Si}_3\text{N}_4$ , forces are always repulsive with the exception of a range around pH 6–7. This means that the tip and sample always bear a charge of the same sign and that the PZC is around pH 6.5. For  $\text{Si}_3\text{N}_4$  on mica the total tip–sample force changes from attractive (for  $\text{pH} \leq 6$ ) to repulsive ( $\text{pH} \cong 8$ ). This is consistent with the fact that mica is negatively charged and the PZC of silicon nitride tip is around  $\text{pH}=6.5$ . For  $\text{Si}_3\text{N}_4$  tips on  $\text{Al}_2\text{O}_3$  the forces are always repulsive, with a repulsion minimum at  $\text{pH}=4.3$ . Hence the PZC is between  $\text{pH}=4$  and  $\text{pH}=5$ . In addition to the experimental results, the authors introduce a simulation method based on the computation of the double-layer force by means of Eq. (6.73), in which the tip and sample charges are calculated according to the site binding theory. Such a method predicts well the qualitative behavior of double-layer force for all the three samples.

Similar results have been obtained by Karaman et al. [142] for an  $\text{Al}_2\text{O}_3$  substrate with aluminum or  $\text{Si}_3\text{N}_4$  tips, and by Senden and Drummond [124] for mica with a silicon nitride tip. In this work, the pH dependence of double-layer force with and without background electrolyte is investigated, and measurements at constant pH with different electrolyte concentrations have been performed. The curves are fitted taking into account both the Van der Waals force and double-layer force (see also Ref. [143]).

Arai et al. [144] have measured the PZC of  $\text{Al}_2\text{O}_3$ ,  $\text{SnO}_2$ , and  $\text{SiO}_2$  with a  $\text{Si}_3\text{N}_4$  tip in buffer. The PZC are measured by monitoring the amplitude of the repulsive or attractive forces at different pH (2–12) at a distance corresponding to  $1/2K_D$ , i.e., 15 nm, in order to make Van der Waals force negligible. For  $\text{Al}_2\text{O}_3$  the results agree with those of Raiteri et al. For  $\text{SnO}_2$  forces are repulsive below  $\text{pH}=6$  and above  $\text{pH}=10$ . The first IEP is due to the PZC of  $\text{Si}_3\text{N}_4$ , the second to that of  $\text{SnO}_2$ . The  $\text{SiO}_2$  has a PZC at  $\text{pH} \cong 2$ , and hence there is only an IEP due to  $\text{Si}_3\text{N}_4$  (see also Refs. [17, 144, 145]). Furthermore, the

authors have studied the pH dependence of forces on a  $\text{SiO}_2$  substrate partially covered with hydrocarbons. They show that both attractive (at  $\text{pH}=3$ ) and repulsive (at  $\text{pH}=10$ ) forces are stronger on bare  $\text{SiO}_2$ . Such a behavior indicates that the surface charge density of bare  $\text{SiO}_2$  is larger than that of the HC-covered quartz. When the quartz is uncharged ( $\text{pH}=6.8$ ), the force–displacement curves on the two substrates overlap and show a slight repulsive double-layer force due to the tip charging.

Lin et al. [137] have once more studied the IEP for a silicon nitride tip on  $\text{SiO}_2$  with two different methods. The IEP due to silicon nitride is  $6.2 \pm 0.2$  and  $5.8 \pm 0.4$  (see also Ref. [62] for the determination of the IEP of  $\text{ZnS}$  and [146] for the determination of the IEP of  $\text{TiO}_2$ ).

Larson et al. [91] have compared the determination of the potentials by means of the AFM with the electrophoretic determination. The system employed was a  $\text{TiO}_2$  colloidal sphere on  $\text{TiO}_2$ . A good agreement between these two methods is found. In later works, the same comparison is established for a silica colloidal sphere on  $\text{TiO}_2$  and on silica [147] and for colloidal silica probes on silica [148]. Once again there is a good agreement between the two methods (see also Ref. [149]).

Butt [150] has acquired force–displacement curves on purple membrane deposited onto alumina at different pH. The purple membrane is negatively charged above  $\text{pH}=3$ , alumina is negatively (positively) charged above (below)  $\text{pH}=8.1$  and the oxidized  $\text{Si}_3\text{N}_4$  tip is always negatively charged. Hence, force–displacement curves are repulsive both on alumina and on purple membrane for pH values larger than 8, while they become attractive on alumina and stay repulsive on purple membrane for  $\text{pH}<8$ . The force is weaker on alumina than on purple membrane. The author could calculate the surface charge density of purple membrane from the ratio of repulsive forces at  $\text{pH}$  10.3 and 9.3 and from the known surface charge density of alumina.

The double-layer force between a very sharp  $\text{ZnO}$  whisker crystal ( $R \cong 10$  nm) and an amino-silanized coverslip has been measured by Aoki et al. [12]. The double-layer force is measured in deionized water, in 0.1 mM and 1 mM  $\text{KCl}$ , and in 0.1 mM and 1 mM  $\text{MgCl}_2$ . The authors take advantage of extremely flexible cantilevers (see Section 1.3) in order to acquire subpiconewton details of the force in the range 0–5 nm.

Drummond and Senden [69] have demonstrated that the double-layer force between a silicon nitride tip and mica in hexadecyltrimethylammonium bromide (CTAB) is intermediate between the theoretical constant surface charge fit and the theoretical constant potential fit (see also Refs. [62, 90, 147, 151, 152]).

Biggs et al. [92, 153] have measured force–distance curves between a gold colloidal sphere and a gold flat substrate in different solutions. They verified the increasing of the strength and decay length of double-layer force with decreasing ion concentrations ( $\text{NaCl}$ , trisodium citrate or gold chloride). By fitting the experimental data both with the constant charge force and the constant potential force, the authors were able to establish that the charge, and hence the adsorbed ions concentration, is constant. Also the dependence of double-layer force on pH was studied (see also Ref. [154]).

Hu and Bard [155] have exploited the double-layer force to monitor the adsorption of sodium dodecyl sulfate (SDS) on functionalized gold substrates. SDS is a negatively charged molecule. When force–distance curves are collected between a silica colloidal probe and silica in SDS solutions at pH 5, the force is always repulsive at any SDS concentration. When the substrate is positively charged gold, the force is attractive for SDS concentrations smaller than  $5 \times 10^{-6}$  M and becomes repulsive for larger concentrations, indicating that the SDS has adsorbed on gold and the surface has become negatively charged. The same behavior is verified for hydrophobic gold (charge reversal for  $10^{-5}$  M SDS concentration). At a certain concentration, called critical micelle concentration (cmc), micelles begin to

form and adsorb onto the substrate.  $\text{Na}^+$  ions begin to bind to micelles in order to compensate the high negative charge of micelles. These bound  $\text{Na}^+$  ions no longer contribute to the diffuse double-layer. By determining the Debye length of the solution, the cmc can be measured in rather good agreement with other methods.

Similar experiments have been performed by Manne et al. [156] with cetyltrimethylammonium bromide (CTAB) on graphite, by Larson et al. [157] with pyridine on gold, by Hu et al. [158] with hexameta-phosphate on CdS, and by Kane and Mulvaney [154] with lead on gold surfaces covered with thiols.

Raiteri et al. [159, 160] have acquired force–displacement curves between a silicon nitride tip and a platinum or gold sample mounted on an electrode. An electric potential was applied between the sample and a counter electrode. Measurements have been performed at pH 9–10 in order to have a negative tip potential. The authors show the existence of three regimes. At negative potentials ( $-0.5$ – $0$  V), as expected, the tip–sample force is dominated by the double-layer force. In this interval of potential the double-layer force is almost constant. Between 0 and 1 V the double-layer force steeply declines. At high potentials ( $> 1$  V) an attractive exponential force onsets, with a decay length of about 50 nm for platinum. The third region is absent in the case of gold (as a matter of fact, anions bind strongly to gold and hence larger positive potentials are needed to make the attractive force appear). In the case of gold, the oxidation of the surface can be revealed due to a sharp increase in adhesion. The calculation of forces by means of the Poisson–Boltzmann theory enables the prediction of the behavior of the system in the first and second region, but the attractive exponential force in the third region cannot be explained with this theory. Force–displacement curves with an applied potential have been measured also by Ishino et al. [161] with similar results.

Hillier et al. [162] have measured the forces between a silica colloidal probe and a gold electrode as a function of the applied potential. They show that the force is strongly repulsive for voltages between  $-0.2$  and  $-0.7$  V, and the double-layer repulsion increases with the decreasing potential. The force becomes slightly repulsive for  $V = -0.1$  V and is attractive for  $V = 0$  V or  $V = 0.1$  V. The authors show that the theoretical curves with constant charge boundary conditions fit better to the experimental data than the constant potential model. The PZC of the system is also determined by plotting the force as a function of the applied voltage at different distances from the sample surface. The authors, both theoretically and experimentally, demonstrate that the point of zero force, i.e., the potential at which the force becomes zero depends on the distance from the sample surface at which forces are measured (going from  $-0.2$  V for forces measured at 25 nm up to  $-0.1$  V for forces measured at 5 nm). The dependence of the PZC on the dimensions of the ions in solutions is also verified by acquiring force–distance curves in NaF, KCl, KBr and KI solutions.

The double-layer force has been measured also at the liquid/liquid interface. Mulvaney et al. [133] have acquired force–distance curves between a glass colloidal particle and an oil droplet in pure water or in a 1 mM sodium dodecyl sulfate (SDS) and 10 mM NaCl solution. In pure water the curve is attractive (see Section 6.3), but when salts are added, they adsorb onto the oil droplet and form a charged monolayer giving rise to the repulsive double-layer force. The measured surface potential of the droplet is less than  $-40$  mV, and the measured Debye length is 10 nm, rather than  $-120$  mV and 0.3 nm, as expected from theory. The authors attribute these large discrepancies to several factors. First of all the contact line between the colloidal sphere and the droplet is affected by strong oscillations depending on the scan rate, probably due to fluctuations in the fluidity of the oil/water interface. Hence, it is very difficult to find the contact distance. Furthermore, since the droplet surface is oscillating, the

contact line might not correspond to a plane of surface charge. Finally, the surfactant SDS is likely to diffuse or migrate out of the region of double-layer interaction so that the surface charge is fluctuating in an undetermined fashion and the oil can flow into the diffuse double-layer.

Drummond and Senden [69] have exploited the double-layer force in order to determine the effective radius  $R_{\text{eff}}$  of a pyramidal tip. The method is based on comparing the forces measured with the colloidal probe and the forces measured with the pyramidal tip at a certain reference distance from the sample surface, thus obtaining a scaling factor that can be used to normalize the force measured at all other distances. The reference distance is chosen in order that the contribution of the Van der Waals force and of other forces is negligible (typically it is 15 or 20 nm) and the entire force is due solely to the double-layer. Furthermore, at such distances, also the difference between the two boundary conditions, i.e., constant charge and constant potential, is negligible.

The study of the double-layer force has perhaps given the best results in AFM studies of forces. The capability of the instrument in characterizing this force with high resolution in any kind of liquid solution and with any kind of surfaces is fully demonstrated. The dependence of the force on the pH of the solution or on the salt concentration has been extensively studied for several kinds of systems. The AFM is routinely used to measure the Debye length of solutions, to determine the PZC of materials, and to calculate the surface charge density of substrates. The forces at a liquid/liquid interface, the dependence of forces on applied potentials and the study of ions adsorption at interfaces are now the most intriguing issues.

### 6.3.3. Dependence on tip shape

As for all other forces, the knowledge of the tip shape is critical for the quantitative measurement of the double-layer force. Arai and Fujihira [163] have calculated the double-layer force according to Eq. (6.55) for a conical sphere-ended tip with the cone angle  $2\theta$  ranging from  $1^\circ$  to  $150^\circ$  and the radius of curvature  $R$  ranging from 1 to 3500 nm and for a spherical tip of the same  $R$ . The calculations were performed in  $10^{-3}$  and  $10^{-4}$  M 1:1 electrolytes. In the  $10^{-3}$  M electrolyte, for  $R = 20$  nm, the forces differ significantly for  $2\theta = 90^\circ$  and  $2\theta = 30^\circ$  (at 10 nm the force is 0.018 nN for the  $90^\circ$  tip and 0.01 nN for the  $30^\circ$  tip). The force for the spherical tip is similar to the force for the  $30^\circ$  tip. Such differences disappear for  $R = 200$  nm. In the  $10^{-4}$  M electrolyte, for  $R = 20$  nm, all the three tips ( $30^\circ$ ,  $90^\circ$ , and spherical) probe different forces (at 1 nm the force is 0.0075 nN for the spherical tip, 0.015 nN for the  $30^\circ$  tip, and 0.0475 nN for the  $90^\circ$  tip). A slight difference persists for the  $30^\circ$  tip and the  $90^\circ$  tip or the spherical tip in the case of  $R = 200$  nm. The difference disappears for  $R = 500$  nm. The authors have also fitted experimental data testing the validity of the low potentials approximation [Eq. (6.70)] for a silicon nitride tip with  $R = 20$  nm and  $2\theta = 30^\circ$  onto a silicon sample in pH 11.4. Experimental data are best fitted by the curve calculated without such an approximation.

Sader et al. [164] have used the linearized Poisson–Boltzmann theory to calculate the electrical double-layer force between identical spheres and have compared the exact value of the force with that obtained with the Derjaguin approximation. Drummond and Senden [69] have utilized this procedure to study the applicability of the Derjaguin approximation. They calculated the ratio of the Derjaguin approximation to the exact interaction between a sphere and a flat surface, assuming that such a ratio for the sphere–plate system is half of that for the sphere–sphere system. They found that, for the constant potential interaction, this ratio is always 1 for  $K_D R > 2$  and is between 1.1 and 1 for  $K_D R < 2$ .

The situation is worst for the constant charge interaction, since the ratio is between 1.5 and 1.05 for  $K_D R < 10$ .

Also microasperities on the tip are of fundamental importance for the quantitative measurement of double-layer force. First of all, surface roughness may lead to a wrong determination of the contact distance and of the position of the plane of charge on the tip and on the sample. One simplistic approach to this problem is to create a macroscopic reference plane at a negative separation equal to half the maximum roughness (average plane of charge) [17, 69]. If the contact occurs at the microasperities separation rather than at the average plane of charge, the derived value of the surface potential is an underestimate [149].

Toikka et al. [62] have showed that, because of microasperities, the liquid is not completely squeezed out at contact and a thin layer of liquid stays between the tip and the sample, thus exerting a double-layer force that decreases the adhesion force. Therefore, the apparent adhesion force depends on pH. The authors found an evidence of this phenomenon measuring adhesion forces at different pH between an iron oxide colloidal probe and silica. The silica surface is negatively charged at  $\text{pH} > 2$ , while the iron oxide surface is negatively charged at  $\text{pH} > 6$ . Thus, for  $\text{pH} > 6$ , there is a repulsive double-layer force. If the liquid was completely squeezed out at contact, adhesion force would not depend on pH. But this is not the case, experimentally, as the adhesion decreases with increasing pH for  $\text{pH} < 6$  and is zero for  $\text{pH} > 6$ .

#### 6.4. Solvation forces

The continuum theories of Van der Waals force and double-layer force cannot describe the mutual interaction of two surfaces approaching at distances smaller than few nanometers, because (a) they are not valid at small separations; and (b) other forces arise. The theory modeling the interactions of the surfaces taking into account only Van der Waals and double-layer forces is called DLVO (Derjaguin–Landau–Verwey–Overbeek) theory [165, 166] and all other surface forces are called non-DLVO forces. The non-DLVO forces can be roughly grouped into four categories: solvation forces, repulsive hydration forces, hydrophobic attractive forces, and steric forces, which are reviewed in the forthcoming sections.

##### 6.4.1. Theory

Solvation forces arise whenever liquid molecules are compelled to order in almost discrete layers between two surfaces. In most cases, these forces exhibit an oscillatory behavior. Consider the way liquid molecules order at an isolated surface or between two surfaces, as indicated schematically in Fig. 30. Computer simulations reveal three different behaviors of liquid molecules at interfaces. At a vapor–liquid or a liquid–liquid interface there is no oscillatory behavior in the density of the liquid, as shown in panel (A) of Fig. 30. At a solid–liquid interface liquid molecules begin to order in layers [167, 168], as shown in panel (B) of Fig. 30. Between two solid surfaces such an ordering increases, as shown in panel (C) of Fig. 30. Even in the absence of any liquid–solid interaction, geometric considerations alone compel the molecules to rearrange in order to fit between the solid surfaces. The variation of this ordering with the separation  $D$  generates the solvation force.



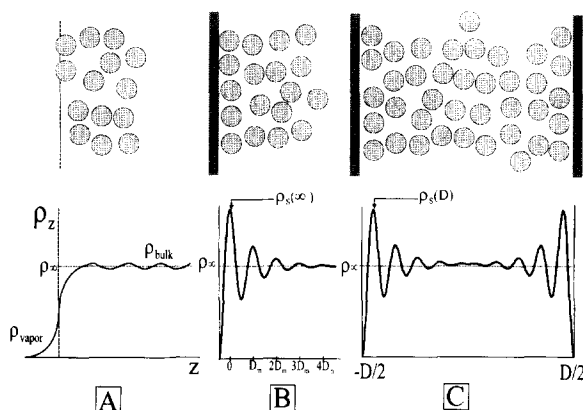


Fig. 30. Schematic representation for the ordering of liquid molecules. Panel (A): vapor–liquid interface. Panel (B): solid–liquid interface. Panel (C): between two solid surfaces. The density of the liquid as a function of position is shown.  $\rho_z$  is the density of the liquid as a function of the distance from the interface,  $\rho_{\text{vapor}}$  is the density of the liquid molecules in the gas phase,  $\rho_{\infty}$  is the bulk density,  $\rho_s(\infty)$  is the density of the liquid at the interface in the case of a single solid–liquid interface,  $\rho_s(D)$  is the density of the liquid at the interface when a second interface is at a distance  $D$ , and  $D_m$  is the radius of the molecule (adapted from [15]).

Use of Eq. (6.56) permits the quantitative description of solvation forces. This equation can be applied to any kind of particle. Therefore, a solvation force occurs whenever there is a variation of density.

Like the density profile, the solvation pressure is an oscillatory function of distance with a period that approximately equals the molecular diameter  $D_m$  and with a range of few molecular diameters. At very small separations, when the last liquid layer has been squeezed out,  $\lim_{D \rightarrow 0} \rho_s = 0$  and the pressure tends to a finite value, deduced from Eq. (6.56):

$$\lim_{D \rightarrow 0} P = -k_B T \rho_s^{\infty}. \quad (6.76)$$

Hence the contact force is adhesive.

In a first approximation the solvation oscillatory force can be described by a cosine exponentially decaying in the form [15]:

$$P(D) \cong -k_B T \rho_s(\infty) \cos \frac{2\pi D}{D_m} e^{-D/D_m}. \quad (6.77)$$

Integrating Eq. (6.77) and inserting  $\rho_s(\infty) = \sqrt{2}/D_m^3$ , yields the work of adhesion due to the solvation force:

$$W(0) \cong \frac{\sqrt{2} k_B T}{4\pi^2 D_m^2}. \quad (6.78)$$

Comparing this value to the work of adhesion due to the Van der Waals force, we conclude that the solvation force dominates the adhesion if the Hamaker constant is less than  $0.2 k_B T$ , i.e., if  $A < 10^{-21}$  J.

O'Shea et al., using the Derjaguin approximation, have calculated the force on a parabolic tip by integrating Eq. (6.77) [169]:

$$F(D) = -k_B T \rho_s(\infty) \frac{2\pi R \sigma}{\sqrt{4\pi^2 + 1}} \cos \frac{2\pi D}{D_m} e^{-D/D_m}, \quad (6.79)$$

in which  $R$  is the radius of curvature of the tip. Kralchevsky and Denkov [170] have obtained a semi-empirical expression of the solvation force as a function of  $D$ ,  $D_m$ , and the molecule volume fraction  $\varphi_V$ . Their expression is

$$P(D) \cong -P_0 \cos \frac{2\pi D}{D_m} \exp \left( \frac{D_m^3}{D_{m1}^2 D_{m2}} - \frac{D}{D_m} \right). \quad (6.80)$$

$P_0$  is given by

$$P_0 = \rho k_B T \frac{1 + \varphi_V + \varphi_V^2 - \varphi_V^3}{(1 - \varphi_V)^3} \quad (6.81)$$

in which  $\varphi_V$  is the molecule volume fraction and  $\rho = 6\varphi_V/\pi D_m^3$ . The dimensionless period  $D_{m1}$  and the dimensionless decay length  $D_{m2}$  are given by

$$\frac{D_{m1}}{D_m} = \sqrt{\frac{2}{3}} + a_1 \Delta\varphi_V + a_2 (\Delta\varphi_V)^2, \quad (6.82)$$

and

$$\frac{D_{m2}}{D_m} = \frac{b_1}{\Delta\varphi_V} - b_2,$$

in which  $\Delta\varphi_V = \pi/3\sqrt{2} - \varphi_V$ ,  $a_1 = 0.23728$ ,  $a_2 = 0.633$ ,  $b_1 = 0.48663$  and  $b_2 = 0.42032$ .

Gelb and Lynden-Bell [171] have studied oscillatory solvation forces between a non-structured tip and a sample in a liquid by means of MD simulations and closed analytical solutions. They showed that the amplitude of the oscillations is linear with the tip radius, with the exception of the first oscillation below  $2D_m$ , which increases non-linearly. The number of oscillations per curve is not sensitive to tip dimension. The oscillations amplitude increases also with decreasing temperature. Neither the temperature nor tip dimensions alter the wavelength of the oscillating force. The density of the liquid has an effect both on oscillation amplitude and wavelength. With increasing density, the amplitude increases and the wavelength decreases. The authors suggested that the dependence of the first oscillatory peak on tip dimension could be exploited in order to calibrate the apex of the tip.

Subsequently, Patrick and Lynden-Bell [172] investigated the effect of structured tips and substrates, by considering an atom-ended tip and a four atoms-ended tip. They showed that, in the case of the “flat” tip, oscillations are about sixfold larger than with the atomic sharp tip. This is due to the fact that the apex atom of the sharp tip compensates the repulsive forces exerted by liquid molecules on the tip. They also showed that the sharp tip is able to trap a liquid molecule into the hole of a fourfold hollow site, thus yielding an additional peak below  $D_m$ . This additional peak is absent when the tip is located above an atop site of the substrate. The trapped liquid molecule stays between the tip and sample for 10–100 ps. Hence, the authors state the incapability of AFM to detect such a phenomenon.

### 6.4.2. Experimental results

The first measurement of solvation forces with an AFM has been performed by O'Shea et al. [173] in octamethylcyclotetrasiloxane (OMCTS) and dodecanol on graphite with a  $\text{Si}_3\text{N}_4$  tip. Both curves are shown in Fig. 31. In OMCTS, there is an oscillatory force superimposed to the attractive Van der Waals force. The oscillatory force results in a series of repulsive surges followed by sudden jumps. The repulsive surge is due to the tip pushing on the structured liquid layer. When the layer beneath the tip is

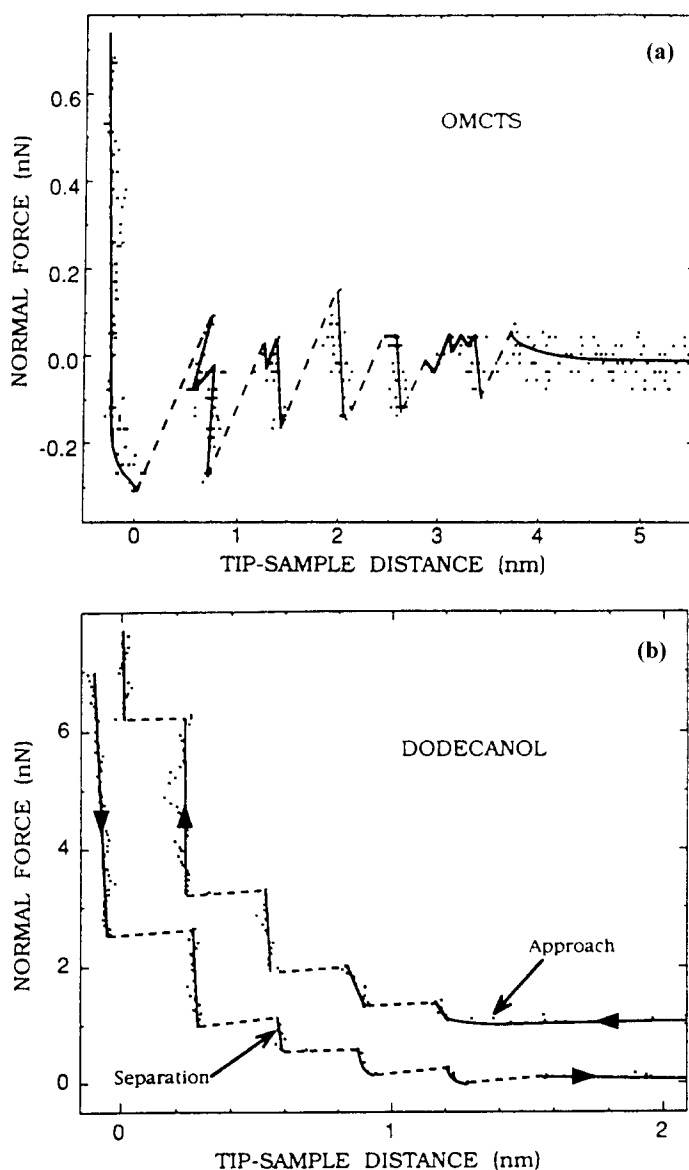


Fig. 31. Force–distance curves in octamethylcyclotetrasiloxane (OMCTS) (panel (a)) and in dodecanol (panel (b)) (reprinted with permission from [173]).

squeezed out, the tip jumps onto the next layer. The curves shown in Fig. 31 seem to have been rearranged by calculating the true tip–sample distance as explained in Section 5.1, but this is not plainly stated in the paper. The oscillations increase with decreasing distance, from  $\cong 0.2$  to  $\cong 0.4$  nN. The period is  $0.73 \pm 0.06$  nm, whereas the molecule has a diameter of 0.8 nm. The small discrepancy between the diameter of the molecule and the period of the oscillations might be due to packing geometry or to errors in the calibration of distances.

In dodecanol, the oscillatory force is superimposed to the repulsive Van der Waals force. It can be observed both on the approach curve and on the withdrawal curve. The amplitude of the oscillations increases with decreasing distance from  $\cong 0.5$  up to  $\cong 2.5$  nN. The period is  $0.35 \pm 0.04$  nm.

The same measurement in OMCTS has been repeated by Cappella et al. [128] on graphite. The authors were not able to measure solvation forces on mica since a few molecules of water (that are likely to remain on a cleaved hydrophilic mica substrate) are likely to destroy the ordering of OMCTS molecules.

Subsequently, O'Shea et al. [169] have acquired both static and dynamic force–distance curves, making the cantilever vibrate by means of an external magnetic field. They were able to detect solvation forces in OMCTS and dodecanol in both modes.

O'Shea et al. could not measure solvation forces in water. This has been done by Hoh et al. [174]. Force–displacement curves were acquired on a glass surface with a silicon nitride tip. The pH was adjusted between 8 and 9 in order to compensate the Van der Waals attractive force with the double-layer repulsive force. A small jump-off-contact is obtained. Recording the jumps-off-contact for several curves, discrete jump-off forces can be observed and the histogram of these jumps shows a well-defined peak at  $1.2 \times 10^{-11}$  N. Also a second peak can be seen, but higher multiples are lost because of peak broadening. The authors explain what they call “quantized adhesion” in two ways. Either the peaks are due to individual hydrogen bonds being resolved, or to solvation forces caused by the ordering of water layers between the tip and sample. A later experiment [175] seems to confirm this latter hypothesis.

### 6.5. Hydration forces

The hydration force probably arises between hydrophilic surfaces, such as silica and mica, since strongly H-bonding surface groups modify the H-bonding network of nearby liquid water molecules. This also occurs when hydrated cations in solution bind to a negatively charged surface at high salt concentration. The hydration force is probably due to the energy needed to dehydrate the ions, as indicated by the fact that the strength and the range of the force increase with the hydration number in the order  $\text{Mg}^{2+} > \text{Ca}^{2+} > \text{Li}^+ \cong \text{Na}^+ > \text{K}^+ > \text{Cs}^+$ . The hydration force is exponential with the distance, but the origin of the exponential decay is still unknown.

Fig. 32 shows force–displacement curves illustrating the effects of the hydration force on mica in  $\text{MgCl}_2$  solutions at different concentrations. At 3 mM  $\text{MgCl}_2$  (first curve), the double-layer repulsive force can be seen. This force disappears at higher concentrations (second and third curve), but at 3M  $\text{MgCl}_2$  the hydration force onsets. The shape of the curve in presence of the hydration force is not very different from that of the curve associated with double-layer force. The presence of the hydration force is inferred from theoretical considerations, since it is known that, for  $[\text{MgCl}_2] > 3$  mM, the double-layer force disappears.

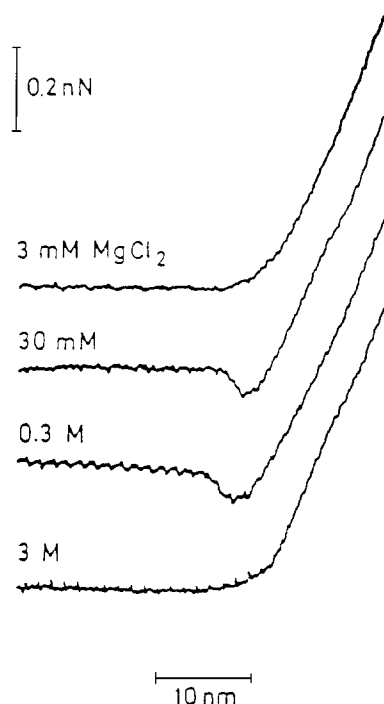


Fig. 32. Force-displacement approach curves illustrating hydration force on mica in  $\text{MgCl}_2$  solutions at different concentrations. Curves are acquired at constant pH with a silicon nitride tip (reprinted with permission from [138]; Biophysical Society).

Similar forces have been observed by Butt [138] for a silicon nitride tip on mica in  $\text{CaCl}_2$  and  $\text{SrCl}_2$ . The force in a 3M  $\text{MgCl}_2$  solution has a decay length of 3 nm and is 0.07 nN strong.

Ducker et al. [17, 176] argue the presence of an hydration force between a silica colloidal sphere and a silica flat plate from the fact that, at small distances ( $< 2\text{--}3$  nm), the measured force does not follow the DLVO fit. The authors list other possible origins of such a mismatch (presence of a gel layer on the silica surface, shift in the position of the plane of surface charge). The authors have observed that the force increases with pH, i.e., with the charge on the surface.

Meagher [139] has measured hydration forces between a silica colloidal sphere and a silicon sample in 0.01 M  $\text{CaCl}_2$  solutions at pH 4.1. He has shown that the theoretical curve based on DLVO theory fits the experimental data well for distances larger than 2 nm. Below 2 nm, however, a repulsive force is probed and no jump-to-contact occurs. The hydration force disappears at pH 5.3 and 1 M or at pH 10.3 and 0.01 M  $\text{CaCl}_2$ . These results suggest that the force is due to the specific adsorption of  $\text{Ca}^{2+}$  onto the silica surface. When the adsorption increases (because pH or ions concentration increase), the force becomes attractive and is well predicted by the DLVO theory.

Atkins and Ninham [151] have measured forces between silica surfaces in 1,2-ethanediol and water. In both solvents, for  $D \geq 3$  nm, the DLVO theory can fit the data, but at smaller separations, a non-DLVO force is measured. This force can be represented by the equation

$$F = \eta R \exp(-\tau D). \quad (6.83)$$

The decay length is  $\tau = 0.97 \pm 0.12$  nm for 1,2-ethanediol and  $\tau = 0.41$  nm for pure water (in good agreement with other measurements). From the fact that the hydration force scales with the radius of the solvent molecules, the authors argue that the hydration force is actually an oscillatory solvation force whose oscillations are smeared out by the roughness of the tip and sample.

Biggs et al. [93] have observed hydration-like forces between gold surfaces in NaCl. Approaching the tip to the sample, an additional repulsion is observed at all concentrations. This repulsion extends over more than 10 nm in distance. The authors attribute this force to the presence of highly hydrated sodium ions near the surface, but do not discuss the differences with the case of mica surfaces.

Larson et al. [147] have observed the hydration force between silica surfaces and silicon on  $\text{TiO}_2$  at high pH. The authors were not able to distinguish between real hydration forces and effects due to tip or sample asperities.

Karaman et al. [142] have studied the forces between an alumina surface and an aluminum or silicon nitride tip. Hydration forces are observed below 3 nm. The cause of this non-DLVO force is attributed to the formation of a thin hydrolyzed oxide gel layer near the surface. This supposition is supported by the fact that, when force–distance curves are repeatedly acquired, initially attractive forces become more and more repulsive, reflecting the characteristic time of formation of the gel layer.

Despite of the importance that hydration forces seem to have in surface science and in biology, this force is still unknown. A lot of theoretical and experimental work is needed. The study of hydration forces by AFM presents several problems. This force acts on a very small range and is in most cases very weak. Often the presence of this force is inferred from the fact that the force–distance curve cannot be fitted by DLVO equations. Errors in the determination of distances, of the Hamaker constants, of the surface charge densities or of surface potentials are likely to affect the fit and to give wrong results. Therefore, this force is rather difficult to characterize quantitatively, and depending on the system, it is not possible to single out the contribution of hydration force to the total tip–sample force.

### 6.6. *Hydrophobic force*

Attractive forces between hydrophobic macroscopic bodies in water have been measured for different systems. The origin of these forces is not well understood, although recently several hypothesis have been proposed.

1. The hydrophobic force could originate from changes of the water structure in the thin layer between hydrophobic surfaces compared to the structure of bulk water [177, 178].
2. The hydrophobic force could be the capillary force due to cavitation in the vicinity of hydrophobic surfaces [179].
3. In another hypothesis, hydrodynamic fluctuations at a hydrophobic surface/water interface are believed to correlate with those at the neighboring interface to give rise to an attractive force [180].
4. Some researchers suggest that hydrophobic force arises from correlated dipole–dipole or dipole–charge interactions [181–183].
5. Others believe that such a force is the result of dipole interactions associated with the large domains of ordered hydrocarbon chains [79, 184, 185].

Since no theory is able to explain the experimental results, data are fitted by an empirical force law in the following form [80]:

$$\frac{F}{R} = C_1 \exp\left(-\frac{D}{D_1}\right) + C_2 \exp\left(-\frac{D}{D_2}\right), \quad (6.84)$$

in which  $C_1$  and  $C_2$  represent the strength of the short- and long-range force respectively, and  $D_1$  and  $D_2$  are the respective decay lengths. The term “short-range” refers to the hydrophobic attraction with decay length of 1–2 nm.

The hydrophobic force can also be fitted by a power law:

$$\frac{F}{R} = -\frac{\Gamma}{6D^2}. \quad (6.85)$$

This equation can fit both short- and long-range hydrophobic force, but there are uncertainties about the exponent, since an exponent slightly larger than 2 can fit experimental data better [79, 184].

Rabinovich and Yoon [80] have measured the hydrophobic force between a silica plate and a glass colloidal probe hydrophobized with octadecyltrichlorosilane (ODTCS). They could measure the parameters in Eqs. (6.84) and (6.85) by measuring the jump-to-contact of the curves, according to the following relationships:

$$(C_{1,2})_{\text{jtc}} = -\frac{k_c D_{1,2}}{R} \exp\left(\frac{D_{\text{jtc}}}{D_{1,2}}\right), \quad \text{and} \quad (\Gamma)_{\text{jtc}} = \frac{3k_c D_{\text{jtc}}^3}{R}, \quad (6.86)$$

in which the subscript “jtc” means that these values are obtained with the “jump” method.

The parameters  $C_1$ ,  $C_2$ , and  $\Gamma$  also have been calculated by fitting the whole curve. The authors show that the “jump” method gives less precise values than the fitting of the whole curve.

Values determined by both methods are correlated with advancing and receding contact angles  $\theta_a$  and  $\theta_r$ . For  $\theta_a = 88^\circ$ , three ranges of the force are evidenced. For  $\theta_a = 95^\circ$ , the force has a short- and a long-range regime, while for  $\theta_a = 102^\circ$ ,  $105^\circ$ , and  $115^\circ$  only the long-range regime can be observed. The parameter  $C_2$  correlates well with the advancing angle (it is smaller than 5 mN/m for  $\theta_a < 95^\circ$  and increases sharply up to 255 mN/m for  $\theta_a > 95^\circ$ ), while  $D_2$  is rather insensitive to  $\theta_a$ . The authors suggest that  $C_2$  depends on the density and hence on the ordering of hydrocarbon chains, while  $D_2$  varies with the distance between the charge or dipole domains. Increasing the density (and the contact angle) results in increasing the ordering, but does not increase the distance between domains.

If  $D_2$  depends on the distance between charge or dipole domains, then it should depend on the Debye length. Following the theory of correlation between charges in lattice arrays on two interacting surfaces, the authors derived the following relation for  $D_2$ :

$$1/D_2 = \sqrt{K_D^2 + (\pi/d_c)^2} \Rightarrow 1/D_2^2 = 10.5z^2[X] + (\pi/d_c)^2, \quad (6.87)$$

in which  $d_c$  is the average charge or dipole distance, and  $z$  and  $[X]$  are the valence and the concentration of the electrolyte. The fit of several experimental data with Eq. (6.87) is not good at high electrolyte concentrations and some experimental data are independent of  $[X]$  (see also Ref. [186]).

The role of gas dissolved in the liquid also was investigated. No difference can be observed between curves acquired in degassed water and water in equilibrium with surrounding air, but this could be due to dissolution of air in water during the measurement. Force–displacement curves in water saturated

with argon (whose solubility is about twice that of nitrogen) show a considerably higher attractive force, and this could be explained with the formation of microbubbles. Gas microbubbles in the liquid gap act in analogous way like liquid bubbles in an air gap, i.e., they contact the tip and sample surface and exert a meniscus force. The sharp increase of the force for  $\theta_a > 95^\circ$  may support this view.

Also the adhesion force is measured. Typical values of  $F_{ad}/R$  are in the range 100–400 mN/m, and the jump-off-contact distance is in the range 20–100 nm.

Forces were measured also for an asymmetric system, i.e., a hydrophobic glass colloidal sphere on a clean hydrophilic silica plate. Only after subtracting DLVO forces, can a short-range hydrophobic force be revealed. In a later work, the same authors have confirmed these results with trimethylchlorosilane [79].

Tsao et al. [185] have studied the dependence of hydrophobic force on the temperature and on the hydrocarbon chain length by collecting force–displacement curves between a silicon nitride tip and mica in several surfactants. All the possible combinations were investigated, i.e., a bare tip and hydrophobic mica, a hydrophobic tip on bare mica, and both tip and mica coated with hydrophobic layers. The surfactants employed are in order of length: dihexadecyl-dimethyl-ammonium (DHDA), dioctadecyl-dimethyl-ammonium (DODA), dieicosyl-dimethyl-ammonium (DEDA), and didoeicosyl-dimethyl-ammonium (DDDA).

The authors show that, in water and at room temperature, the force is weaker for the shortest chain (DHDA), and that the force is stronger for the hydrophobic mica–bare tip system than for the bare mica–hydrophobic tip system, whereas for the symmetric situation the force is much weaker. These results are in contrast with the results of Rabinovich and Yoon, and these two latter authors believe that Tsao et al. have not accounted for DLVO forces.

When raising the temperature, both DHDA and DODA show only the Van der Waals attraction, while the force for DEDA is substantially the same. This result correlates well with the chain transition temperatures in water.

Hydrophobic forces are present also in ethylene glycol. In this solvent, the force is the same for DDDA and DEDA, but smaller for DODA. As a matter of fact, the melting temperature of DODA in ethylene glycol is comparable with the melting point of DHDA in water (34°C and 35°C). The fact that hydrophobic forces have been observed in solvents different from water indicates that this force cannot be attributed to the unique structural properties of water. Further, the force is associated with the solid phase of the surfactant, and disappears as the melting point is approached. No long-range attractive hydrophobic forces were observed when a bare  $\text{Si}_3\text{N}_4$  tip approached either a Teflon or polyethylene substrate in either water or ethylene glycol, indicating that these two polymeric surfaces do not possess specific molecular structures that give rise to the hydrophobic force. Finally, the force depends on salt concentration.

Similar results were obtained in another work [187] measuring forces for DHDA, DEDA and DODA at 25°C, 40°C, and 50°C. The first derivative of the force can be fitted with Eq. (6.84). All parameters (i.e., both the force strength and the force decay length) decrease with increasing temperature for any surfactant. The correlation between the force and the chain order is studied by means of molecularly resolved AFM images.

Rabinovich et al. [184] have correlated the results of Tsao et al. with the order parameter  $S$  of the hydrocarbon chain defined as  $S = (3\langle \cos^2 \gamma \rangle - 1)/2$ , where  $\gamma$  is the angle between the axis normal to the surface and the molecular axis. When  $S = 1$ , all molecules are parallel to each other and perpendicular to the surface, whereas when the molecules are randomly oriented,  $S = 0$ . They showed that the



parameters  $C_1$ ,  $C_2$ ,  $D_1$  and  $D_2$  found by Tsao et al. have a linear dependence on  $S$  and that  $C_2$ ,  $D_1$  and  $D_2$  are proportional to  $S$ , while for  $C_1$  there is a critical parameter  $S_{cr} \cong 0.3$ , below which the short range hydrophobic force ceases to exist. Also the parameter  $\Gamma$  shows a linear dependence on  $S$  for  $S > 0.3$ , and is zero for  $S < 0.3$ . The authors explain such a dependence as arising from the changes in water structure induced by the ordering of the hydrocarbon chains. The changes in water structure may in turn create dipole domains which are comparable in size with the observed decay length.

Forces also differ for a single-chain surfactant and a double-chain surfactant, in that single-chain surfactants have only the short-range term. This effect can be explained with ordering considerations. On the one hand, since the adsorption is controlled mainly by the Coulomb attraction between the polar head and the negative charge site on the mica substrate, the area occupied by each single-chain molecule is actually the area per charge site ( $\cong 0.5 \text{ nm}^2$ ). This area is approximately twice as large as the cross-sectional area of the hydrocarbon chain. As a result, the hydrocarbon chain are disordered and exhibit only the short-range hydrophobic force. On the other hand, the double-chain molecules give rise to large values of  $S$  and consequently to long-range forces.

Mantel et al. [188] have measured forces between a silanated glass colloidal particle and stainless steel. They found that hydrophobic forces are absent in the non-treated stainless steel that contains a certain amount of hydrophilic potassium salt (2%). When the stainless steel is treated with Ar plasma and then rinsed with water or simply rinsed with water, the potassium decreases and the hydrophobic force onsets. The authors argue that, when hydrophilic and hydrophobic groups are present in close proximity on a surface, the attractive hydrophobic force and the repulsive hydration force, both arising from changes in water structure, are not simply additive.

Several studies have addressed the interaction of an AFM tip with bubbles or droplets in water. Even if they are not directly related to hydrophobic force, these results are reviewed in this section, since some evidences exist that hydrophobic interactions could also arise from microbubbles in water.

Butt [189] has acquired force–displacement curves and force–time curves between a glass colloidal particle and an air bubble in water. When the glass particle is hydrophilic, force–displacement curves show a repulsive region when the tip and the bubble are in contact. Several oscillations are observed on such contact lines. These oscillations are reproducible when the force is kept below a certain threshold value and become hysteretic if the force exceeds the threshold value, because the particle snaps into the bubble. When the glass particle is hydrophobic the curve is completely different. A strong attractive force makes the tip snap into the bubble and a very high force is needed to pull it out of the bubble. Looking at the time-resolved analysis of the signal, an oscillatory behavior of the acceleration can be noted. The attractive force is calculated between 580 and 1000 nN. Also the forces between a glass particle and a water droplet in air are considered. The system hydrophilic tip–water droplet is analogous to the hydrophobic glass sphere–air bubble. The tip snaps in at a certain distance with a force between 230 and 406 nN, but the time-resolved analysis reveals that the process is 2–4 times faster, probably because of the lower effective mass of the cantilever in air. Also the hydrophobic glass snaps into the water droplet. The author has offered no explanation for this result.

Ducker et al. [190] have acquired force–displacement curves between a silica colloidal particle and an air bubble in water. When the silica particle was hydrophobic their results are in agreement with those of Butt. Also in the case of an hydrophilic silica particle an attractive force of about  $-0.2 \text{ mN/m}$  is detected. After the tip has jumped into contact, the curves show a contact line from which the stiffness of the air bubble can be measured ( $0.065 \pm 0.005 \text{ N/m}$ ). This contact line is probably due to the presence of a thin wetting film between the silica particle and the air bubble and the repulsive force is

associated with the high energy needed to squeeze it out. The attractive force cannot be accounted for by the sum of the force needed to squeeze out the wetting film and of the DLVO force. The sum of the jump distances due to these two forces is 26 nm, while the measured jump distance is 44 nm. By measuring jump-in distances at different salt concentrations, the authors ascertained that the force was not electrostatic. When sodium dodecylsulfate (SDS) is adsorbed at the air–water interface, forces become repulsive and are in good agreement with DLVO theory at all distances, with deviations only below 2 nm. The deviations are now repulsive. The measured forces are interpreted on the basis of capillary bridging forces. When the tip is pushed against the sample, the air bubble “wets” the tip surface and exerts an attractive force, even when the tip is hydrophobic. Also hydrophobic forces are interpreted as arising from the air microbubbles nucleation between the two surfaces.

### 6.7. Specific forces

Specific forces are non-covalent forces that generate very strong adhesion between molecular groups. The term “specific force” is somewhat ambiguous, but can refer to all forces that are present only between a specific pair of molecules. Most of the interactions between biological molecules are due to specific forces. The specific bond between a couple of molecules arises from the cooperation of several non-covalent bonds (e.g., hydrogen or ionic bonds) favored by the geometry of the interacting molecules. The two molecules may fit together by means of a “lock and key” mechanism. Specific bonds, like covalent bonds, have a precise stoichiometry.

In order to measure specific forces with the AFM, it is necessary to functionalize the tips by covering them with one of the two molecules under study. A typical withdrawal curve between a tip functionalized with odorant binding protein (OBP) and a sample functionalized with its antibody is shown in Fig. 33. At a distance of 7 nm from the contact point, there is a jump-off-contact of 230 pN due to

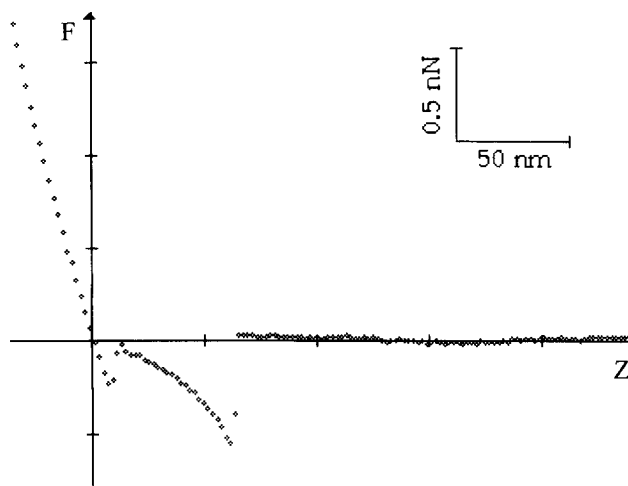


Fig. 33. Withdrawal force–displacement curve with specific adhesion between odorant binding protein (OBP) and anti-OBP. After the jump-off-contact, there is a non-linear attractive surge due to the stretching of a group of OBP molecules, that suddenly breaks at a distance of 65 nm.

aspecific adhesion (mainly Van der Waals force). After the jump-off-contact, the cantilever returns to the rest position (zero deflection). The tip and sample seem to be detached, but some OBP and antibody molecules are still bound by bonds which are stretched along the subsequent attractive surge until the force reaches a rupture value. The discontinuity of 550 pN at a distance of 65 nm marks the detachment due to the rupture of the specific bonds between the OBP and its antibody. Since the complementary molecules may be located at different points on the tip and on the sample, they may detach at different distances, and hence a specific force–displacement curve may present multiple detachments. The rupture occurs at a distance of 65 nm because the sample undergoes important deformations and because a whole group of OBP molecules is likely to be stretched.

In principle, the jump-off-contact and the specific detachment may overlap so that the specific force is hidden by the aspecific (Van der Waals) adhesion. In some works this problem is solved by acquiring curves in a particular liquid that has a small Van der Waals force. In some others, one or both of the interacting molecules are attached to the tip and/or to the sample via a spacer. The spacer is a molecule few nanometers long which is stretched during the detachment and makes the rupture of the bond occur at distances greater than those of the aspecific detachment. The curve in Fig. 33 is acquired with the antibody bound via a spacer.

In general, four different parameters can directly be extracted from specific force–displacement curves:

- (1) the force of a single binding event;
- (2) the rupture distance depending on the length of the molecules involved and on how they are immobilized;
- (3) the adhesion probability, i.e., the ratio between the number of force–displacement curves with specific detachment and the total number of force–displacement curves, depending on surface coverage and experimental conditions; and
- (4) the elasticity of the stretching region.

The first study on specific forces with the AFM is that of Florin et al. [20], who measured the forces between avidin and biotin molecules. They used  $\text{Si}_3\text{N}_4$  commercial tips functionalized with avidin on agarose beads functionalized with biotin. The measured adhesion between the two surfaces is of about 20 nN. When blocking the surfaces with an excess of free avidin or biotin, the adhesion diminishes down to  $\cong 200$  pN. The force–displacement curve displays both the jump-off-contact of about 1 nN and the specific detachment occurring at different distances. The forces do not diminish in presence of excess bovine serum albumin (BSA), which proves the specificity of the interactions. Histograms of the measured rupture forces show several regularly spaced peaks, that correspond to multiples of an elementary force of  $160 \pm 20$  pN. This force is attributed to the interaction of a single avidin–biotin couple. In similar experiments carried out with quite hard substrates (silicon oxide instead of agarose beads) no quantized adhesion is found. A likely explanation is that the soft agarose bead allows the molecules to arrange parallel to each other, so that the force does not depend on the relative positions. This is not possible with hard substrates. As a matter of fact, the spacer is often employed also to make the molecules more mobile and to allow such a rearrangement. Measurements have been repeated with desthiobiotin and iminobiotin, leading to the identification of an elementary force of  $125 \pm 20$  and  $85 \pm 15$  pN, respectively.

In a later work [191], the same authors have showed that there is a meaningful correlation between the measured elementary forces and the enthalpy of the system.

Lee et al. [192] have measured the forces acting in the following three systems: biotin–biotin, biotin–streptavidin blocked with biotin, and biotin–streptavidin. The force between biotin and biotin is characterized by a repulsive steric or hydration force and no adhesion is present. Between biotin and blocked streptavidin there is an adhesion force of  $60 \pm 40$  pN. This adhesion force may be due either to the increase of the Hamaker constant or to the decrease of the repulsive force. The interaction between biotin and streptavidin is dominated by an adhesion of  $340 \pm 120$  pN. The authors note that, because of steric limitations due to the immobilization of molecules on the surfaces, not all the biotin molecules within the contact area (nearly 10 molecules) bind a streptavidin molecule. They observe that no specific forces can be measured with a sharper tip.

The same kind of measurements have been performed by Allen et al. [193]. No adhesion is observed when the streptavidin is blocked, and the specific force turns out to be  $409 \pm 166$  pN.

Bowen et al. [194] have measured the forces between a tip and a sample both functionalized with bovine serum albumin (BSA). The forces are studied at different NaCl concentrations and at different pH.

Lee et al. [195] have measured specific interactions between DNA strands. All experiments are performed in a 0.1 M NaCl solution. In a first experiment the forces between two complementary strands, namely (ACTG)<sub>5</sub> and (CAGT)<sub>5</sub>, are measured. The adhesion forces are found to be grouped in four categories. One of these ( $0.48 \pm 0.1$  nN) corresponds to the non-specific force between non-complementary strands, the other three correspond to interactions between DNA oligonucleotides of 12 ( $0.83 \pm 0.11$  nN), 16 ( $1.11 \pm 0.13$ ) or 20 ( $1.52 \pm 0.19$ ) base pairs. Since the oligonucleotides have different lengths, the bond breaks at different distances and forces. In a second experiment both the tip and the sample are functionalized with cytosine. On the tip, inosine polymers of 160 base pairs are bound to the cytosine. In withdrawal curves, retarded detachments at distances of about 240 nm can be observed with specific forces between 1 and 2 nN. These detachments are due to the bond between the inosine and the cytosine on the sample, and 240 nm is the length of the inosine polymer. Forces between couples of cytosine polymers also can be measured.

Boland and Ratner [196] have measured the hydrogen bonding between DNA complementary pairs (adenine–thymine and cytosine–guanine). The histograms of the attractive forces show in both cases quantized forces with an elementary force of 54 pN. The attractive and adhesive force turn out to be the sum of 12 and about 60 base pair interactions, respectively. Several control experiments with non-complementary base pairs and bare surfaces are performed.

Noy et al. [197] have measured the forces between two complementary DNA 14-mers fixed on a gold surface via HS(CH<sub>2</sub>)<sub>16</sub> spacers. The force–displacement curve between DNA strands has the typical shape of the curve in Fig. 33. The specific adhesion is of 0.45 nN (one couple) or of 0.9 nN (two couples) and the detachment occurs at a distance of about 18 nm. When employing two non-complementary DNA strands, force–displacement curves show a jump-off-contact of about 100 pN at zero separation. The authors have also studied the stretching of the DNA duplex formed before the detachment. They observe that the part of the curve between the origin and the specific detachment, i.e., the part resulting from the stretching of the duplex, is characterized by three regions: a short elastic region, a pronounced flat region where the separation increases rapidly under almost constant force (indicative of a structural transformation), and a relatively stiff elastic region immediately prior to duplex separation. The elastic constant of the DNA duplex before and after the structural transformation is determined. The qualitative and quantitative results obtained by Noy et al. are in agreement with some theoretical predictions.

Nakagawa et al. [198] have measured the forces between trichlorosilanes of different length. The experiments are performed in ethanol in order to decrease the Van der Waals adhesion. The force is found to be proportional to the length of the chain for  $n \geq 8$ , whereas there is no specific force for  $n < 8$ , since the chains are randomly oriented and do not bind to each other.

Dammer et al. [199] have measured the specific forces between biotin and anti-biotin molecules. The adhesion force between the two functionalized surfaces is between 100 and 200 pN, while it is almost zero for the non-functionalized tip. The authors carry out several control experiments (non-functionalized tip, blocked tip or sample, low and high pH, and non-specific antibody) in order to ensure the measured forces to be specific. In these control experiments the measured adhesion is between 20% and 40% of the specific adhesion. The most interesting control is the one at low and high pH. When the pH of the solution is changed, the ionic and hydrogen bonds cooperating in the specific bond are affected, and the specific interaction may disappear. By blocking the biotin on the tip, it is possible to decrease the interaction force, and the histogram of the force values shows four peaks that are multiples of an elementary force of  $60 \pm 10$  pN.

Antigen–antibody specific forces also have been measured by Hinterdorfer et al. [200] taking advantage of the complex human serum albumin (HSA)–anti-HSA. The authors solve the problem of distinguishing aspecific and specific adhesion by introducing a spacer (polyethylene glycol PEG, 8 nm long) between the tip and/or the sample and the molecules. The elementary force is about 240 pN and the rupture distance goes up to 30 nm, which is the sum of the lengths of one HSA molecule, one anti-HSA molecule and two PEG molecules. Both the values of the rupture force and distance indicate that individual molecule–molecule bonds have been measured. The distribution of the rupture lengths is bimodal. The two different distances correspond to the detachment of the two arms of the asymmetric antibody molecule.

Stuart and Hlady [201] have measured the forces between fluorescein and anti-fluorescyl IgG. They have observed that the strength of the specific force and the overall shape of the curve depend on the contact time between the tip and sample. When the tip and sample are allowed to stay in contact only for 2 s, no adhesion is observed. When the contact time is increased up to 1 or 2 min, the adhesion increases up to nearly 4 nN and the jump-off-contact distance is 200 nm (1 min) or 400 nm (2 min).

Allen et al. [202] have directly monitored specific interactions between ferritin and anti-ferritin in an immunoassay system. The measured forces are quantized with steps of  $49 \pm 10$  pN. This elementary force is attributed to individual unbinding events.

Vinckier et al. [203] have measured specific forces between the chaperonin GroEL and several proteins (synthase, lactamase, bovine serum albumin, and peroxidase). Forces are measured with and without ATP in solution and both with native and denatured proteins. The authors have verified that, for all proteins, the forces are greater in the denatured form than in the native form. The forces in presence of ATP are greatly reduced, since the ATP modifies the GroEL structure. This happens both for native and denatured proteins. The non-hydrolyzable ATP analog, the ATP $\gamma$ S, does not affect the forces. The authors have shown that the specific interaction depends on the hydrophobicity of the tip. Hydrophilic functionalized tips do not interact with GroEL, while hydrophobic functionalized tips show the same interaction as tips functionalized with proteins. Finally the authors have studied the dependence of the measured force on the tip diameter. The force between GroEL and lactamase goes from 110 pN for a tip diameter of 35 nm up to 380 pN for a tip diameter of 160 nm.

Nakajima et al. [204] have studied the actin–heavy meromyosin (HMM) interaction, with a single HMM molecule bound onto the tip surface via a long chain made up of biotin, ultravidin, and

acrylamide. The force–displacement curves look like the curve in Fig. 33, but sometimes the jump-off-contact overlaps with the specific detachment. The histogram of the pull-off forces is bimodal with peaks at 14.8 and 24.7 pN. The two forces correspond to the binding of only one head or of both heads.

Dammer et al. [205] have measured the forces between proteoglycans which are involved in cell adhesion (AP). The pull-off is once again retarded due to the lifting and extending of string-like arms. In 10 mM  $\text{Ca}^{2+}$  the mean force is 125 pN and the probability of binding is 60%. The force and the probability are greatly reduced in 2 mM  $\text{Ca}^{2+}$  and 2 mM  $\text{Ca}^{2+}$  + 10 mM  $\text{Mg}^{2+}$  (40 pN and 12%), showing that the AFM measurement is able to assess the  $\text{Ca}^{2+}$ -selectivity of an interaction. Also control experiments with a blocking antibody have been performed. Subsequently, Fritz et al. [206] have studied the elasticity of the proteoglycans and have addressed the extension dependence of the force. These authors have found that the extension dependence of the force is non-linear and that two fundamental phenomena contribute to this non-linear response. At low force regimes (below some 10 pN) the response is dominated by entropic effects and the cantilever works against a random thermal motion, which tends to curl the molecule. At higher forces, the elasticity of single-chain segments has to be considered. The data of Fritz et al. are best fitted by an empirical formula which considers only the entropic elasticity [207]. If  $X$  is the extension of the molecule, the force due to entropic elasticity is described by

$$F(X) = \frac{k_B T}{\Xi} \left[ \frac{1}{4} \left( 1 - \frac{X}{\Sigma} \right)^{-2} - \frac{1}{4} + \frac{X}{\Sigma} \right], \quad (6.88)$$

in which  $\Xi$  is the persistence length and  $\Sigma$  is the contour length. If the chain molecule is represented as a chain of  $N$  freely joint rigid rods, the persistence length is defined as the length of each rod and the contour length as  $N\Xi$ .

Kikuchi et al. [208] have studied the stretching of polystyrene chains and have fitted the extension–force law with the extended Langevin equation:

$$X(F) = \left[ \coth \left( \frac{lF}{k_B T} \right) - \frac{k_B T}{lF} \right] \left( n_m l + \frac{n_m F}{k_m} \right), \quad (7.89)$$

in which  $X$  is the extension,  $F$  the force,  $l$  the length of the monomer,  $n_m$  the number of monomers, and  $k_m$  is the elastic constant of the monomer.

Bowen et al. [209] have functionalized the tip with a cell, namely a yeast cell (*Saccharomyces cerevisiae*). They have measured the adhesion of the cell onto a mica surface. After a first jump-off-contact, there is a long stretching region (about 200 nm) in which the attractive force increases slowly until the cell detaches from the surface. Such curves are contact time dependent. If the contact time is increased to 5 min, then the adhesion increases and the stretching region becomes longer (about 500 nm).

Another group of experiments deals with the determination of the single bond force for other kind of forces (mainly Van der Waals and hydrogen bond). In these experiments the tip is functionalized with a non-biological molecule. So, strictly speaking, the measured forces are not specific. This kind of AFM measurements is often called “chemical force microscopy”, since the functionalized tip behaves as a chemical sensor.

Van der Vegte and Hadziannou [210] have measured the forces between the chemical groups listed in Table 8. All measurements have been performed in ethanol. Once the adhesion force is measured, the

Table 8

Measured mean adhesion in nN and single bond force in pN (bold) between the listed functional groups as determined in Ref. [210]. All measurements have been performed between functionalized tips and samples in ethanol.

tip	Substrate									
	CH <sub>3</sub>		OH		NH <sub>2</sub>		COOH		CONH <sub>2</sub>	
CH <sub>3</sub>	0.9	<b>81</b>	0.3	<b>57</b>	0.3	<b>59</b>	0.3	<b>61</b>	0.3	<b>60</b>
OH	0.1	<b>50</b>	0.9	<b>101</b>	1.2	<b>113</b>	1.2	<b>112</b>	1.4	<b>117</b>
NH <sub>2</sub>	0.2	<b>54</b>	0.5	<b>88</b>	0.8	<b>98</b>	0.7	<b>95</b>	0.9	<b>100</b>
COOH	0.7	<b>95</b>	1.2	<b>109</b>	1.2	<b>105</b>	1.3	<b>114</b>	2.2	<b>137</b>
CONH <sub>2</sub>	0.3	<b>62</b>	1.3	<b>110</b>	1.2	<b>102</b>	2.2	<b>125</b>	1.8	<b>120</b>

JKR theory is applied in order to determine the surface energy  $\gamma$  and the single bond force. The number of interacting molecules is given by the ratio between the contact area, determined by means of JKR theory, and the characteristic area of the molecules. The elementary force is the total adhesion divided by the number of molecules. The trend expected on the basis of the kind of interactions governing the adhesion is qualitatively observed. In the case of a CH<sub>3</sub>-modified tip, for which only the Van der Waals interaction plays a role, small adhesive forces are found. The cohesive CH<sub>3</sub>–CH<sub>3</sub> Van der Waals interaction is larger than the Van der Waals force for dissimilar pairs, e.g., CH<sub>3</sub>–COOH. As regards OH, NH<sub>2</sub>, COOH, or CONH<sub>2</sub> tips, again the pure Van der Waals interaction with a CH<sub>3</sub> surface gives the lowest adhesive force. The measured hydrogen bond forces are in qualitative agreement with hydrogen bond energies.

Williams et al. [211] have developed a method to calculate the single bond forces. This method is based on Poisson statistics. The distribution of adhesion forces in an AFM experiment can be modeled with the Poisson statistics provided the two following assumptions are satisfied:

- (1) the total adhesion force is the sum of a finite number of discrete, independent chemical bonds;
- (2) these bonds form randomly, and have similar force values. Accordingly, if the single bond force is  $F$ , the mean force  $\bar{F}$  and the variance  $\sigma_F^2$  are given by

$$\bar{F} = mF \quad \text{and} \quad \sigma_F^2 = mF^2, \quad (6.90)$$

in which  $m$  is the mean number of events, i.e., of bonds. The single bond force can be calculated as

$$F = \frac{\sigma_F^2}{\bar{F}}. \quad (6.91)$$

If an additional aspecific force  $F_0$  is present, then

$$\bar{F} = mF + F_0 \quad \text{and} \quad \sigma_F^2 = mF^2 = F\bar{F} - FF_0. \quad (6.92)$$

Thus a plot of  $\sigma_F^2$  versus  $\bar{F}$  gives both the single bond force (the slope of the linear plot) and the aspecific force (the intercept of the line is  $-FF_0$ ). The authors determine the single bond force for the Van der Waals force and for the hydrogen bond by measuring forces between a Si<sub>3</sub>N<sub>4</sub> tip and gold or mica, respectively. Measurements are performed in water. Different values of the mean force are obtained with different tips of randomly varying  $R$ . The Van der Waals single bond force is  $60 \pm 3$  pN

and the hydrogen bond force is  $181 \pm 35$  pN. Van der Vegte and Hadziannou have compared their results with the one of Williams et al. verifying that the two are in good agreement. The same result for the Van der Waals single bond force had been previously found by Han et al. [212]. In this latter work, Van der Waals and hydrogen single bond forces have been measured in several liquids. The elementary Van der Waals force between an  $\text{Si}_3\text{N}_4$  tip and a gold surface is  $59.7 \pm 4.2$  pN in water and  $29.2 \pm 4.2$  pN in propanol. As regards the COOH-functionalized tip and sample, only the measurement in hexane gives good results for the hydrogen bond elementary force. As a matter of fact, by analyzing the distance dependence of the force, the authors find a power law of  $-3$  (i.e., the expected value for hydrogen bond interaction) only in hexane. The elementary force is  $70.6 \pm 2.4$  pN. In water and propanol the interaction is repulsive, and is due to the double-layer force.

Subsequently, the hydrogen bond elementary force has been measured by functionalizing the tip and sample with SH- or OH-terminated silanes [213] or  $\text{CH}_2\text{Br}$  or  $\text{CH}_3$  terminated silanes [214]. In the first experiment forces are measured in water,  $(\text{CH}_3)_2\text{SO}$ , and propanol. The OH–OH single bond force in water is  $119 \pm 16$  pN. The single bond force is linearly dependent on the dielectric constant of the medium and the value in vacuum is extrapolated to  $362 \pm 10$  pN for OH–OH and to  $136 \pm 6$  pN for SH–SH. In the second experiment forces are measured in water, methanol, and propanol. The hydrogen bond elementary force between unmodified  $\text{Si}_3\text{N}_4$  tips and unmodified  $\text{SiO}_2$  substrates in water is  $236 \pm 20$  pN. The Van der Waals single bond force between bromo-undecyl-trimethoxy-silanes is  $31 \pm 5$  pN in water,  $75 \pm 11$  pN in methanol, and  $101 \pm 3$  pN in propanol. Finally, the force between octa-decyl-trimethoxy-silanes has a single bond force of  $658 \pm 68$  pN in water and of  $281 \pm 35$  pN in propanol. The authors have suggested that these rather high values of force are due to a hydrophobic interaction.

Thomas et al. [215] have measured the elementary forces for the systems  $\text{CH}_3\text{--CH}_3$ ,  $\text{NH}_2\text{--NH}_2$ ,  $\text{COOH--COOH}$ , and  $\text{NH}_2\text{--COOH}$ , taking advantage of their force feedback controlled AFM (see Section 3.2). Measurements are performed in an evacuated chamber successively filled with dry  $\text{N}_2$ . The force feedback method allows the authors to sample the whole curve and hence to work in dry  $\text{N}_2$ . The work of adhesion and the single bond energy are determined on the basis of DMT theory. As in the work of van der Vegte and Hadziannou, the  $\text{CH}_3\text{--CH}_3$  interaction is a pure Van der Waals interaction and the calculated surface energy ( $30 \text{ mJ/m}^2$ ) is in agreement with Lifshitz theory. All other interactions are interpreted as hydrogen bond interactions.

Van der Vegte and Hadziannou [216] have studied the pH dependence of the adhesion between a tip and a sample both functionalized with the same acidic or basic group. The curves depicting such a dependence resemble the solution titration curves for acids. Furthermore, the authors have deduced the degree of ionization and the effective pK from force–displacement curves. For COOH and  $\text{PO}_3\text{H}_2$  the adhesion decreases for  $\text{pH} > 4$  and reaches zero for  $\text{pH} = 10$ . The initially strong adhesion is due to hydrogen bonds, whereas the final zero adhesion is due to the double-layer force that increases as more and more COOH or  $\text{PO}_3\text{H}_2$  groups become deprotonated. In the case of the diprotic acid  $\text{PO}_3\text{H}_2$  the two steps in the ionization, i.e.,  $\text{PO}_3\text{H}_2 \rightarrow \text{PO}_3\text{H}^-$  and  $\text{PO}_3\text{H}^- \rightarrow \text{PO}_3^{2-}$ , are clearly observed. In the case of  $\text{NH}_2$ , the adhesion is zero for  $\text{pH} < 4$  and then increases with pH. The initial zero adhesion is due to the fact that, at low pH, the  $\text{NH}_2$  groups are positively charged ( $\text{NH}_3^+$ ) and give rise to a double-layer repulsion. Finally, the adhesion is pH-independent for the  $\text{CH}_3$  groups.

All the experiments dealing with the determination of specific forces have demonstrated the importance of the techniques based on functionalized tips. Functionalized tips have a chemical selectivity that greatly enhances the capabilities of the AFM in studying biological interactions and



mechanisms of molecular recognition. Specific forces, single bond forces, the stretching of molecules and acid–base interactions have been extensively investigated. A lot of work is still necessary, however, in order to clarify several details of such measurements and to make the technique reliable and simple to perform.

### 6.8. Steric, depletion, and bridge forces

Steric forces arise when the interacting surfaces are spatially diffuse, i.e., they have thermally mobile surface groups. A common type of thermally diffuse surface is a surface covered with chain molecules which dangle out into the solution, e.g., a polymer-covered surface. When another surface is approached confining these dangling chains in a well-defined volume, a repulsive entropic force known as “steric” or “overlap” force arises.

The extension of a polymer chain into the solution depends on three interactions: the solvent–solvent, the monomer–monomer, and the monomer–solvent interaction. If the solvent is a “ $\theta$ ” solvent, i.e., the polymer–solvent interaction is equivalent to the polymer–polymer interaction, the root mean square radius of the polymer coil  $R_g$  is given by  $R_g = l\sqrt{n_m}/\sqrt{6}$  [217], where  $n_m$  is the number of monomers and  $l$  is the length of a monomer. In real solvents the effective size of the coil is referred to as Flory radius  $R_F$  [217]. In a “good” solvent there is a repulsion between the monomers and  $R_F \cong ln_m^{3/5}$ . In a poor solvent the polymer collapses to a globular shape. The thickness of the polymer layer  $T_L$  depends also on the coverage of surface. If  $N$  is the number of chains per unit area, in a good solvent  $T_L$  is given by [15]:

$$T_L = \sqrt[3]{NR_F^{5/3}}. \quad (6.93)$$

Once two polymer-covered surfaces are closer than  $2T_L$ , there is a repulsive pressure  $P(D)$  given by [218]:

$$P(D) \cong k_B TN^{3/2} [(2T_L/D)^{9/4} - (D/2T_L)^{3/4}], \quad (6.94)$$

in which  $D$  is the distance between the surfaces.

Eq. (6.94) is the expression of the steric repulsion. The first term in Eq. (6.94) is the osmotic repulsion between the coils favoring their stretching, while the second term comes from the elastic energy of the chains acting against the stretching. The energy scales as  $D^{-5/4}$  at short distances and as  $D^{-2}$  at large distances [218]. For  $0.2 < D/2T_L < 0.9$  the pressure is adequately described by [218]:

$$P(D) \cong 100k_B TN^{3/2} \exp\left(-\pi \frac{D}{T_L}\right). \quad (6.95)$$

When the surfaces are closer than  $R_g$  the coils are pushed out of the gap resulting in a reduced polymer concentration between the surfaces. Another kind of force arises, no longer due to the confining of the dangling chains which have now been pushed out. This new force is called the depletion force. If the polymer concentration is  $\rho$ , applying the contact value theorem, Eq. (6.56), yields the attractive depletion free energy per unit area [15]:

$$\lim_{D \rightarrow 0} W = -\rho R_g k_B T \quad (6.96)$$

Fleer et al. [219] have deduced the depletion force in the form:

$$F = \pi \frac{\mu}{V_m} (D + 2R)(D - 2T_L), \quad (6.97)$$

in which  $\mu$  is the chemical potential of the solvent,  $V_m$  the solvent molecular volume, and  $R$  is the radius of the tip. The net interaction between two polymer-covered surfaces depends also on the polymer–surface interactions and on the availability of free binding sites on the opposite surface. If there are free binding sites on the opposite surface, some polymer coils will form bridges between the two surfaces and give rise to an attractive bridging force. The bridging force decays roughly exponentially with distance, with a decay length that is close to  $R_F$ .

Biggs [220] has measured the forces between a surface covered with polyacrylic acid (molecular weight  $M_w = 750\,000$ ,  $R_g = 56\text{ nm}$ ) and a zirconia colloidal probe of radius  $\cong 300\text{ nm}$ . After checking that the forces between the probe and a zirconia plate can be completely described by the DLVO theory, the forces between the probe and polymer-covered surface are measured. At high coverage, the measured force on approach increases approximately exponentially from 180 to 20 nm. From 20 nm down to contact a much steeper force is observed. Note that 180 nm corresponds to  $3.5R_g$ . The same features are observed in the withdrawal curve. However, between 0 and 20 nm, the withdrawal force is always lower than the approaching one. The non-exponential force at  $D < 20\text{ nm}$  is attributed to a severe restriction of the degrees of freedom of the polymer and/or to surface roughness. The interaction energy scales as  $D^{-5/4}$  for  $D < 20\text{ nm}$  and as  $D^{-2}$  for  $D > 20\text{ nm}$ . At lower concentration there is a bridging force between the two surfaces that extends up to  $2\text{ }\mu\text{m}$ .

Similar results have been obtained by Frank and Belfort [221] with extracellular polysaccharides (EPS) of varying molecular weight. Bridging forces are present both for anionic EPS and for anionic dextran and extend up to 1300 and 400 nm, respectively. No adhesion is observed. Bridging forces are absent in the case of neutral EPS. The authors have addressed also the effects of the ionic strength of the solution and have showed that bridging forces disappear in a 0.1 M KCl solution (see also Ref. [222]).

Braithwaite et al. [223] have studied the forces between a glass surface and a glass colloidal particle both covered with polyethylene oxide ( $M_w = 56\,000$ ). The force–distance curves acquired after 35 min of incubation show a weak jump to contact at a distance of 45 nm ( $5R_g$ ) and few bridging events due to some dangling polymers adsorbing to free binding sites onto the tip. The withdrawal curve has a slide-off-contact behavior. If sufficient time elapses (e.g., 24 h), the surface is fully covered by the polymer, and no bridging occurs. Only an exponential repulsion, beginning at 90 nm ( $12R_g$ ), is present. If the time of incubation is very long, e.g., some days, force–distance curves are not reproducible. Their common features are an exponential repulsion beginning at very large distances (e.g., 1600 nm) and an irregular slide-off-contact with multiple stages of detachment. When covering only one of the surfaces, one would expect that the bridging events increase, since in this case one of the surfaces is entirely available for polymer adsorption. Measured curves do not support this supposition. The authors have attributed the discrepancy to the rather high scan rate. Finally, the authors have proved that the interactions between polymer-covered surfaces depend on repeated contacts. If the scan rate is high compared to the relaxation time of the polymer and the exerted loads are important, both attractive and adhesive forces change with repeated contacts. In particular, the adhesion diminishes.

Ducker and Clarke [224] have studied the forces between a silicon nitride tip and a silicon nitride flat surface in solutions containing a zwitterionic surfactant adsorbed on the surfaces. A steric repulsion of

about 2 nN is observed in approach and withdrawal curves from 5.2 down to 3.3 nm. The thickness of the two surfactant bilayers, measured by X-ray diffraction, is 4.6 nm. Repeating the contact on the same point of the sample for several times (300), the adhesion diminishes, and the range of the repulsive force extends a distance approximately equal to an additional bilayer. If the tip is moved to a new position, the shape of the curve returns to the previous one.

Steric and bridge forces depend not only on the repeated contacts, but also on the scan rate. Biggs [225] has acquired force–distance curves between a zirconia colloidal probe ( $R = 2\text{--}3\text{ }\mu\text{m}$ ) and a zirconia plate covered with polyvinyl pyrrolidone (PVP) ( $M_w = 40\,000$ ). The curves show a steric repulsion on approaching and several bridging events on withdrawal. The bridging events occur at rather small distances (40–80 nm) and assume the shape of small attractive surges similar to those obtained with functionalized tips (see Fig. 33). The steric force begins at a distance of 70–80 nm ( $\cong 5R_g$ ). Force–distance curves are acquired at different scan rates (9.8 – 0.1 Hz). Differences between the data at 1 Hz and below are insignificant, but, at higher scan rates, the steric force becomes steeper and steeper as the scan rate increases. These data indicate that there is a limiting collision velocity below which the response of the repulsive potential is rate independent. Also the bridge force depends on the scan rate. The jump-off-contact and the bridging events are present only at scan rates of 1 Hz or lower (see also Ref. [222]).

The weakness of the depletion force relative to the double-layer and the Van der Waals forces has made the measurement of this force particularly difficult. Milling and Biggs [226] have measured the depletion force between a silica colloidal particle ( $R = 3.8\text{ }\mu\text{m}$ ) and a silica plate, both covered with octadecyl alcohol ( $\text{SiO}_2\text{--C}_{18}$ ). The Van der Waals force is minimized by performing the experiment in a solution of polydimethyl siloxane in cyclohexane. The curves show a short-range steric repulsion ( $D = 3\text{--}4\text{ nm}$ ) followed by a strong attractive force, beginning at 20 nm of separation. The experimental data are in good agreement with Eq. (6.97).

Milling and Vincent [227] have studied the depletion force between silica surfaces in solutions of polyacrylic acid (PAA). The depletion force depends on the PAA concentration. In particular, the surface separation at which the depletion force reaches its maximum decreases with increasing PAA concentration. Also the effect of electrolytes ( $\text{NaNO}_3$  and  $\text{NaOH}$ ) is studied. The data are fitted with Eq. (6.97).

## 7. Imaging based on force–distance curves

Force–distance curves also may be acquired in all points throughout the scanned area in order to compare tip–sample interactions at different regions of the sample, i.e., in order to study the spatial variation of interactions.

In such a case, different scanning methods are employed and the quantities of interest (e.g., adhesion force) are drawn from the acquired curves and represented in two-dimensional maps. There are some common drawbacks of these imaging techniques. The acquisition of the force–distance curves on each point of the scanned surface can require some minutes, and almost all the information on the sample physico-chemical properties is drawn from maps obtained by data post-processing. An advantage of the use of force curves is the possibility of doing the lateral movement with the tip away from the surface, and hence the capability to avoid the dragging of the sample.

Two methods of acquisition are commonly used. In one method [228] force–displacement curves are acquired, irrespectively of sample topography, over a given distance range  $\Delta Z$  beginning from a fixed

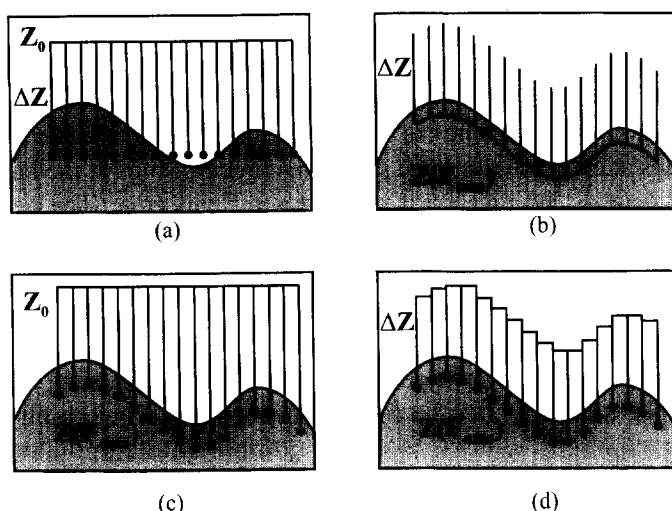


Fig. 34. The four main scanning methods proposed for the sequential acquisition of force–displacement curves on a surface, i.e. Radmacher's method (panel (a)), Baselt and Baldeschwieler method (panel (b)), Koleske et al. method (panel (c)), and Cappella et al. method (panel (d)). In Radmacher's method force–distance curves are acquired, irrespectively of sample topography, over a given distance range  $\Delta Z$  beginning from a fixed distance at  $Z_0$ . In the other three methods, force–distance curves are acquired beginning from a point of maximum load  $Z(F_{\max})$ . In the method of Baselt and Baldeschwieler the tip moves laterally in contact with the sample surface. The sampled distance  $\Delta Z$  is fixed. In the method of Koleske et al. the height at which the tip is retracted is the same for all the points, so that the sampled distance varies for each curve. In Cappella et al. method the scanned volume follows the sample surface.

distance at  $Z_0$  (panel (a), Fig. 34). The scanned volume does not follow the sample surface and there is no control on forces exerted on the cantilever by the sample. Therefore, the cantilever may bend excessively and break on protrusions or fail to contact the surface on depressions. With another method [229–231], force–displacement curves are acquired beginning from a point of maximum load  $Z(F_{\max})$ . Thus the tip does not fail to contact the surface at each point and cantilever damage because of excessive load cannot occur. This method provides another advantage in that all force–displacement curves are related to the same value of force and can be compared directly. After Baselt and Baldeschwieler, the former method is referred to as “absolute mode” and the latter as “relative mode”. In the relative mode different techniques are used to control the force on the cantilever so that it does not exceed the maximum load.

In Baselt and Baldeschwieler's work [229] a controller based on a digital signal processor is used and the tip moves laterally in contact with the sample surface. During this process the feedback loop is active and keeps the cantilever deflection, and hence the tip–sample force, constant. Then, at each point, the feedback is stopped and the tip is first withdrawn and then approached step by step. At each step the deflection is measured and stored. The sampled distance  $\Delta Z$  is fixed, during both the withdrawal and the approaching ramp, as indicated in panel (b) of Fig. 34. In this technique part of the advantage of the force–distance curves method is lost because the tip is moved while it is in contact with the sample and so there is the risk of dragging the sample.

In the method of Koleske et al. [230] the computer stops the approach when the force reaches the maximum load and then the sample is retracted to a given height (see panel (c) in Fig. 34). This height

is the same for all the points so that the sampled distance varies for each curve. Therefore some acquisition time is wasted in sampling points too far from the surface, where there is no interaction.

Yet another method has been proposed by Cappella et al. [231], as illustrated in panel (d) of Fig. 34. At the beginning the tip is away from the surface. Then the sample is approached to the cantilever, acquiring deflection values at each pre-assigned step. This process stops at the height  $Z(F_{\max})$ , when the force exerted on the tip reaches the pre-assigned maximum load  $F_{\max}$ . The deflections in the last  $N_z$  approaching points are stored. Next the sample is retracted step by step for a distance  $\Delta Z$  and the deflections in the next  $N_z$  withdrawal points are acquired and stored. After the acquisition of these  $N_z$  values, the sample is withdrawn for a further short distance in order to insure that the vertical distance between the cantilever and the subsequent  $XY$  point is never less than  $\Delta Z$ . In this method lateral  $XY$  movements are done away from the surface, avoiding sample dragging. The scanned volume follows the sample surface so that no time is spent in acquiring points too far away from the surface.

The spatial variation of surface interactions may be represented in several ways. One of the most effective of these is constructed by plotting all the curves of one scanned line together. Even if the details of the interactions are sometimes lost, this kind of plot has the advantage of immediately comparing all the regions of the curve. Hence, differences in elasticity, attractive forces, and adhesion forces are simultaneously observable. A more refined analysis of the spatial variation of interactions can be performed by extracting several different parameters from the curves. Single parameters can be plotted on a gray scale as a function of the  $XY$  coordinates. In some works, such quantities, e.g., adhesion force, are directly acquired. Finally, a kind of image proposed by Cappella et al., namely the force-slice, has the advantage of showing the differences between different regions of the sample on line, i.e., during the acquisition and not in post-processing images.

### 7.1. Comparative curve plotting

A collection of withdrawal force–displacement curves acquired on a peroxidase deposit on silicon is showed in Fig. 35 [231]. All the curves belong to a scanning line. The peroxidase occupies the center of the scanned area. The curves are acquired in deionized water with an  $\text{Si}_3\text{N}_4$  tip. The first and last curves of the scan line are on clean silicon. These curves have a jump-off-contact of approximately 3 nN. The adhesion on the peroxidase region (central curves) is smaller, allowing the two regions to be distinguished. Similar images have been obtained for a deposit of fluorescein iso-thiocyanate (FITC) on silicon in air [231]. In this case, the differences in adhesion are due to the different degree of hydrophobicity of the two materials.

Radmacher et al. [228] have studied the spatial variation of interactions for a thin nickel film evaporated onto mica, a glass slide partially covered with chromium, and lysozyme adsorbed on mica. All measurements are performed in water. In all the three cases, large differences between the different regions of the sample are shown by plotting the curves of a scan line altogether. The presence of nickel islands on mica and of chromium regions on glass is revealed by higher attractive and adhesive forces, probably due to a hydrophobic meniscus exerted by contaminants adsorbed onto nickel or chromium. For the lysozyme, only the withdrawal curves show clear differences, since the adhesion is weaker on lysozyme than on mica.

Radmacher et al. [232] have mapped force–displacement curves on human platelets showing differences between the substrate and the cell mainly due to elasticity.

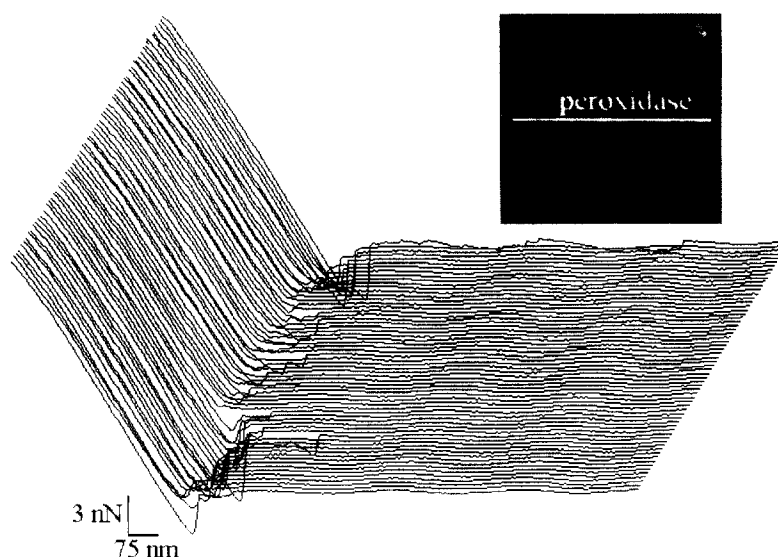


Fig. 35. Line of withdrawal force–displacement curves acquired on a peroxidase deposit on silicon. The inset is a lateral force image showing the area occupied by the peroxidase. The curves on silicon (first and last curves) have a rather large adhesion whereas the adhesion on the peroxidase region (central curves) is smaller, thereby allowing to be distinguished (adapted from [231]).

Rotsch and Radmacher [11] have studied the spatial variation of forces on a vesicle suspension of dimethyl dioctadecyl ammonium bromide (DODAB) on mica in pure water. Several kind of images are presented. The most effective images are the lines of force–displacement curves, in which all the regions of the curves can be immediately compared. In pure water, the DODAB is positively charged, while the mica is negatively charged. The silicon nitride tip is negatively charged. Hence the double-layer force is attractive between DODAB and silicon nitride and repulsive between mica and silicon nitride. Such a difference can be exploited to identify the DODAB vesicles in images collecting both the approach and the withdrawal force–distance curves. The lateral resolution is better than 25 nm.

## 7.2. Force-slices

A force-slice collects the deflections of the cantilever on each point of the scanned area at a given distance from the point of maximum load. The kind of information given by a force-slice depends on the considered distance. If the force-slice collects points on the contact line, then it provides information about the different stiffnesses of the regions of the sample. Force-slices collecting points near the jump-to-contact or the jump-off-contact give information about attractive or adhesive forces. Three force-slices at different distances from the deflection of maximum load (300, 350, and 440 nm) are shown in Fig. 36 for a sample of FITC on silicon. Two squares of FITC are partially evident. The lighter parts of the image indicate the larger value of deflection. In each of the three slice-forces, it is evident that the tip is still adhering to silicon, although it is already detached from the FITC deposit. Hence the points on silicon are black (negative deflections) and the points on FITC are white (zero deflection). The number of points on FITC at which the tip is already detached increases with the distance. Such differences in adhesion are due to capillary forces. Similar images have been presented for a deposit of



Fig. 36. Three force-slices at a distance of 300, 350, and 440 nm from the deflection of maximum load on a FITC deposit on silicon. The lighter parts of the image indicate the larger value of deflection. In each force-slice the tip is still adhering to silicon (black points), while it is already detached from the FITC deposit (white points). The number of points on FITC at which the tip is already detached increases with the distance (adapted from [231]).

peroxidase on silicon in water [228], where differences are mainly due to different Hamaker constants. Meaningful information can be drawn from force-slices only if they are acquired in a relative mode. Otherwise, the curves are not referred to the same maximal value of load and deflections cannot be compared.

Information obtained from force-slices is qualitative. A force-slice does not contain either the value of the pull-off force or the distance at which the tip pulls off the sample. It simply shows that, at a certain distance from sample surface, pull-off has already occurred in some regions of the scanned area, while in other regions it has still to occur.

### 7.3. Mapping parameters drawn from force–distance curves

Three maps obtained in air on a phosphatidyl choline (PC) deposit on silicon are shown in Fig. 37. The maps show the jump-off-contact force (panel (a)), the jump-off-contact distance (panel (b)), and the adhesion (panel (c)). The adhesion is calculated as the area below the zero line in the withdrawal

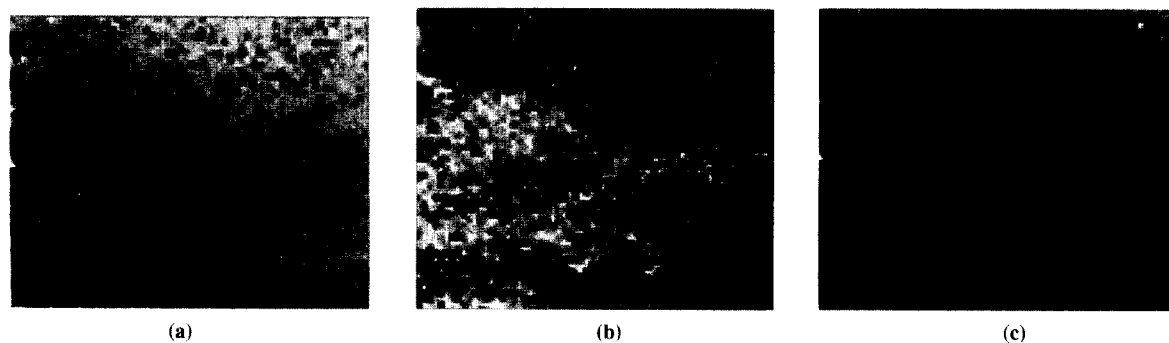


Fig. 37. Three maps obtained in air on a phosphatidylcholine deposit on silicon. The maps show the jump-off-contact (panel (a)), the jump-off-contact distance (panel (b)), and the adhesion (panel (c)).

curve. The PC regions show a smaller pull-off force (darker points), a smaller pull-off distance (lighter points) and a weaker adhesion (darker points). The different quantities are calculated off-line from the curves and mapped. Since curves are acquired in air, the differences between the two regions are due to differences in the thickness of the adsorbed water layer.

The first direct measurement showing the spatial dependence of adhesion is that of Mizes et al. [19]. These authors have performed adhesion studies on a sample of both doped and undoped polycarbonate in air. The adhesion is mainly due to meniscus force. First of all, the authors check that the adhesion is not affected by repeated contacts. Subsequently, they show that the adhesion depends on sample topography and on materials properties. As a matter of fact, three ridges are present on the sample. The adhesion is smaller on the edge of the ridges, since the contact area, and hence the meniscus force, is smaller. On the bottom of the ridge the adhesion is larger, since a larger portion of the tip is in contact with the sample surface, as indicated schematically in Fig. 38. Hence, adhesion maps are able to reveal differences in sample topography. Differences in sample materials can also be detected. The authors demonstrate the capability of adhesion maps to distinguish between doped and undoped polycarbonate. Since measurements are performed in air, the differences detected are probably differences in materials hydrophobicity.

Differences in materials hydrophobicity have been exploited also by Berger et al. [100]. These authors show adhesion maps acquired in air on a di-myristoyl-phosphatidyl-ethanolamine (DMPE) deposit. The thicker domains have a smaller adhesion.

The same kind of measurements has been performed by Sasaki et al. [233] on a Cr grating on glass. The authors show the topography dependence of adhesion and propose a method to draw pull-off and pull-on forces from the curves. Also pull-on maps are presented.

Dufrene et al. [234] have performed similar measurements on a mixed di-stearoyl-phosphatidyl-ethanol-amine (DSPE) and di-oleoyl-phosphatidyl-ethanol-amine (DOPE). Adhesion maps in air show that the microscopic domains of DSPE have a smaller adhesion ( $6.1 \pm 0.2$  nN) than the DOPE matrix

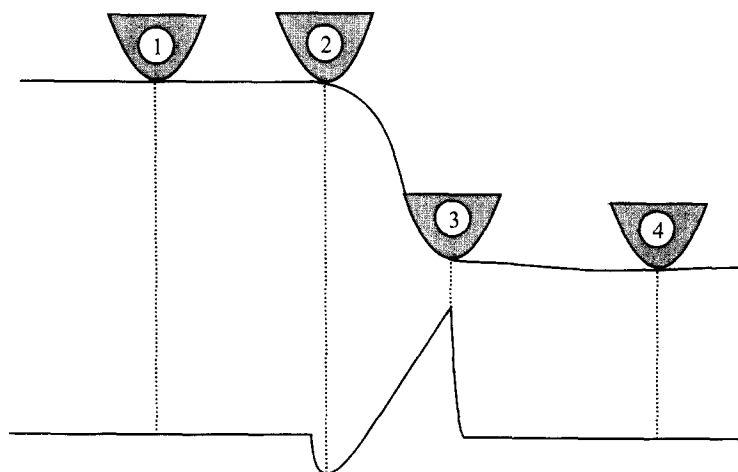


Fig. 38. Schematic representation of the dependence of the adhesion force on sample topography. The top line represents the sample and the tip in four different positions, the bottom line represents the corresponding adhesion. Various tip positions are indicated by the labels (1)–(4). In (2), a smaller fraction of the tip surface comes in contact with the sample surface than in (1) or in (4), whereas the contact area is larger in (3). The adhesion varies accordingly.



( $10.5 \pm 0.2$  nN). Differences in adhesion are attributed to differences in materials elasticity and deformation. Adhesion maps also in water are presented. Once again the adhesion is larger on DOPE. The authors attribute such differences to the presence of repulsive hydration or steric forces acting only between DSPE and the  $\text{Si}_3\text{N}_4$  tip.

Baselt and Baldeschwieler [229] have acquired force–displacement curves in air on collagen fibrils. These authors have used both the relative and the absolute mode of acquisition. Force-slices were considered too difficult to interpret. Hence, different parameters are drawn from the force curves and mapped. The authors have shown adhesion and hardness maps.

Koleske et al. [230] have shown maps of pull-on and pull-off force on a calibration grating in air.

Van der Werf et al. [125] have acquired on-line the pull-off force, the pull-off distance, and the adhesion on a mercaptopentadecane–gold layer on glass. The pull-off force is determined by means of a peak revelator. The pull-off distance and the adhesion are detected by means of two integrators. Several images of different samples obtained in air and in water are presented. The method proposed by the author presents some advantages when compared with usual mapping methods. However, it cannot be effectively used for every kind of sample and interaction.

Ishino et al. [235] have mapped forces on patterns of stearyl-mercaptan, stearic acid, stearyl alcohol, and stearylamine monolayers on  $\text{SiO}_2$ . All the patterns are well imaged in jump-in and jump-off maps. The stearyl-mercaptan is hydrophobic and uncharged in water at pH 5.6, whereas the silicon nitride tip and the bare silicon oxide surface are both negatively charged. In jump-in maps, forces on stearyl-mercaptan are about 1.5 nN greater than on  $\text{SiO}_2$ . The other three patterns are imaged in adhesion maps. In all the three cases, adhesion forces are greater on the monolayers than on the bare silicon oxide. Differences in attractive and adhesion forces are due mainly to the double-layer force (see also Ref. [12]).

#### 7.4. Affinity imaging

The mapping of specific forces by means of functionalized tips has often been called “affinity imaging”, since functionalized tips are able to single out the complementary molecule on a sample.

Hinterdorfer et al. [200] have localized single antigens on samples with a low density of molecules. These authors have performed a scanning of the sample along a given direction. In withdrawal curves, the specific detachments occur only at certain positions. A plot of the number of specific detachments vs. the position gives a gaussian distribution whose peak corresponds to the position of the molecule.

Ludwig et al. [236] have mapped the interaction between avidin and biotin. In their experiment, the tip is functionalized with biotin. The sample is a pattern of avidin lying between squares covered with bovine serum albumin (BSA). The authors show that the usual topographic images are not able to distinguish between the avidin and the BSA regions. On the contrary, the adhesion maps reveal the presence of avidin due to the different interactions. When free biotin is added in solution, the topographic image is always the same, but in the adhesion maps, the contrast disappears.

All the techniques designed to map surface forces have revealed the AFM as a useful tool for distinguishing different materials and for the imaging of materials that would be damaged by other imaging modes. Several technical problems are still unsolved, e.g., the long time needed for the acquisition. Nevertheless these new kinds of “surface spectroscopy” are very promising.

## 8. Synopsis

The AFM has been used to study surface forces and materials elasticity by means of force–distance curves since 1989. Since then, this technique has become a fundamental tool for several kinds of measurements.

In certain cases the technique gives satisfactory results and the interpretation of data no longer presents problems. Thus, AFM force–distance curves are the most used technique for measurements of the elasticity of materials, for the study of the wetting properties of lubricants, for the determination of the Hamaker constants in certain liquids, for the characterization of the double-layer force (determination of the Debye lengths, of the surface charges, and of the PZT of oxide-like materials), and for the study of solvation and hydrophobic forces. New applications are constantly being found. The technique is becoming increasingly reliable in the characterization of the degree of hydrophobicity, in the study of the adsorption of salts and polymer at liquid/solid or liquid/liquid interfaces, and in the characterization of steric and bridging forces.

The most promising results are the recent studies of forces at a liquid/liquid interface, the applications of functionalized tips for the study of specific forces and for the determination of single bond forces, and the development of techniques for the mapping of physico-chemical properties drawn from force–distance curves.

The main problems in the quantitative application of this technique are the uncertainties in the cantilever elastic constant and in the tip shape. Although procedures have been developed to calibrate the elastic constant of the cantilever, some cases, e.g., with functionalized tips, remain in which it is quite difficult to apply such techniques effectively. The main problem in AFM force–distance curve interpretation is, however, the lack of reliable methods to determine the tip shape and dimensions. The use of a big colloidal tip of known dimensions is not an effective solution to this problem, because one of the major advantages of the AFM, namely the capability of probing local interactions on a very little area, is lost. Only a better knowledge of the shape of the tip will increase the reliability of quantitative measurements of forces, above all the Van der Waals and the double-layer forces, and will permit the study of still uncharacterized forces such as hydrophobic or hydration forces.

## Acknowledgements

We thank Dr. C. Ascoli for fruitful discussions and suggestions and Dr. N. Burnham for kindly providing several original figures.

We acknowledge Dr. C.B. Duke, Editor, for carefully proofreading and editing the manuscript and for suggesting a lot of corrections and improvements to our first draft. Despite his gigantic effort, this article remains imperfect.

## List of acronyms

AFM	atomic force microscope
BSA	bovine serum albumin
cmc	critical micelle concentration
CTAB	hexadecyl trimethyl ammonium bromide

DDDA	didoeicosyl dimethyl ammonium
DEDA	dieicosyl dimethyl ammonium
DHDA	dihexadecyl dimethyl ammonium
DLVO	Derjaguin–Landau–Verwey–Overbeek
DMT	Derjaguin–Müller–Toporov theory
DODA	dioctadecyl dimethyl ammonium
DODAB	dimethyl dioctadecyl ammonium bromide
DMPE	dimyristoyl phosphatidyl ethanol amine
DOPE	dioleoyl phosphatidyl ethanol amine
DSPE	distearoyl phosphatidyl ethanol amine
EAM	embedded-atom method
EPS	extracellular polysaccharides
FEP	polytetrafluoroethylene- <i>co</i> -hexafluoropropylene
FITC	fluorescein iso-thiocyanate
HMM	heavy meromyosin
HSA	human serum albumin
IEP	isoelectric point
JKR	Johnson–Kendall–Roberts theory
MD	molecular dynamics
OBP	odorant binding protein
OMCTS	octamethyl cyclotetrasiloxane
PAA	polyacrylic acid
PAN	polyacrylonitrile
PC	phosphatidyl choline
PDMS	polydimethylsiloxane
PEG	polyethylene glycol
PFD	perfluorodecalin
PFMCH	perfluoromethyl cyclohexane
PMMA	poly(methyl methacrylate)
PMN	lead magnesium niobate
PTFE	polytetrafluoroethylene
PVDF	polyvinylidene fluoride
PVP	polyvinyl pyrrolidone
PZC	point of zero charge
PZT	lead zirconate titanate
SDS	sodium dodecyl sulfate
SEM	scanning electron microscope
SFA	surface force apparatus
TEM	transmission electron microscope

### List of symbols

$a$	contact radius
$a_0$	contact radius at zero load

<b>A, B</b>	coefficients of the attractive and repulsive Lennard-Jones terms
$A$	Hamaker constant
$A_r$	retarded Hamaker constant
$A$	amplitude of vibration
$\bar{A}$	dimensionless contact radius
$\alpha, \alpha_{ij}$	electronic polarizability
$\alpha_{\text{exc}}$	excess polarizability
$c$	light speed
$C$	constant of the atom–atom pair potential
$C_1, C_2$	coefficients of the short- and long-range hydrophobic forces
$C_K, C_D, C_L$	Keesom, Debye, and London coefficients
$\gamma$	damping coefficient
$\gamma_L$	surface energy of liquids
$\gamma_S$	surface energy
$\gamma_t, \gamma_s$	tip and sample surface energies
$d_{31}$	proportionality factor between the voltage applied to the piezo and the displacement
$\Delta D, D$	tip–sample distance
$D_1, D_2$	decay lengths of the short- and long-range hydrophobic forces
$D_{\text{jtc}}$	tip–sample distance at which the jump-to-contact occurs
$\Delta\delta_c, \delta_c$	deflection of the cantilever
$(\delta_c)_{\text{jtc}}$	cantilever deflection at which the jump-to-contact occurs
$\delta, \delta_s$	sample deformation
$\bar{\delta}$	dimensionless deformation
$e$	electron charge
$E$	Young modulus
$E$	electric field
$\varepsilon_0$	vacuum dielectric constant
$\varepsilon, \varepsilon_i$	dielectric constant of the medium
$f, F, \Delta F$	force
$F_0 \exp(i\omega t)$	exciting force exerted by the vibrating piezo
$F_{\text{ad}}$	adhesion force
$F_{\text{attr}}$	attractive force
$F_{\infty}^{\text{sp}}, F_0^{\text{sp}}$	Coulomb force between a spherical tip and a flat sample at short and long distances
$F^c$	Coulomb force between a conical tip and a flat sample
$\bar{F}$	mean force
$\bar{F}$	dimensionless force
$\phi$	phase
$\varphi_P$	Pashley parameter
$\varphi_V$	molecule volume fraction
$g$	gravitational acceleration
$h$	Planck constant
$H$	zero load elastic deformation

$H'$	zero load plastic indentation
$H_r$	relative humidity
$k_B$	Boltzmann constant
$k_c$	elastic constant of the cantilever
$k_m$	monomer elastic constant
$K$	reduced Young modulus
$K_D$	Debye length
$l$	monomer length
$L$	length of rectangular cantilevers
$L_1, L_2$	characteristic lengths of “V-shaped” cantilevers
$L_p$	length of the piezotube
$\lambda$	Maugis parameter
$\Lambda$	wavelength of the laser
$m$	mean number of bonds
$m^*$	effective mass of the cantilever
$m_t$	mass of the tip
$M_w$	molecular weight
$\mu$	chemical potential
$\mu_M$	Müller parameter
$n, n_i$	refractive index
$n^*, n_i^*$	low-frequency indices
$n_m$	number of monomers
$N$	number of polymer chains per unit area
$\nu$	Poisson ratio
$\nu_e$	absorption frequency
$\nu_i$	ionization frequency
$\nu_n$	frequency
$\nu_{rot}$	average rotational relaxation frequency
$P$	pressure
$Q_0, Q$	quality factor of the cantilever
$r$	distance
$R$	tip radius
$R_F$	Flory radius
$R_g$	mean square radius of the polymer coil
$R_m$	molecular diameter
$\rho, \rho_i$	number density
$S$	order parameter
$\sigma, \sigma_i$	surface charge density
$\sigma_A$	Attard and Parker parameter
$\sigma_F^2$	variance
$t$	time
$t_c$	thickness of the cantilever
$T$	temperature
$T_L$	thickness of the polymer layer

$\theta_a, \theta_r$	advancing and receding contact angles
$u_i$	dipole moment
$u_M$	penetration depth in the liquid layer
$U_c$	Hooke's elastic potential of the cantilever
$U_{cs}$	tip–sample interaction potential
$U_s$	Hooke's elastic potential of the sample
$U_{tot}$	total tip–sample potential
$v$	vertical scan rate
$\Delta V, V$	voltage
$V_m$	molecular volume
$w$	width of rectangular cantilevers
$w_p$	wall thickness of the piezotube
$w_K, w_D, w_L, w_{vdW}$	Keesom, Debye, London, and total Van der Waals potentials
$W$	adhesion work at contact
$W$	width of V-shaped cantilevers
$X$	extension of the molecule
$\Xi$	persistence length
$\Sigma$	contour length
$y$	distance from the center of the contact circle
$z$	valence
$\Delta Z, Z$	distance between the sample surface and the rest position of the cantilever
$\omega_0, \omega'_0$	resonance frequency of the cantilever
$\Omega$	proportionality factor between $V$ and $\delta_c$
$\psi$	electrostatic potential

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