High-temperature electron localization in dense He gas

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We report accurate measurements of the mobility of excess electrons in high-density helium gas in extended ranges of temperature [$(26 \leqslant T \leqslant 77)$ K] and density [$(0.05 \leqslant N \leqslant 10.0)$ atoms nm^{-3}]. The aim is the investigation of the combined effect of temperature and density on the formation and dynamics of localized electron states. The main result of the experiment is that the formation of localized states essentially depends on the relative balance of fluid dilation energy, repulsive electron-atom interaction energy, and thermal energy. As a consequence, the onset of localization depends on the medium disorder through gas temperature and density. The transition from delocalized to localized states shifts to larger densities as temperature is increased. This behavior can be understood in terms of a simple model of electron self-trapping in a spherically symmetric square well.

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I. INTRODUCTION

The transport properties of excess electrons in dense noble gases and liquids give useful information on the electron states in a disordered medium and on the relationship between the electron-atom interaction and the properties of the fluid. The electron behavior depends on the strength of its coupling with the gas atoms and on the response function of the gas itself. Therefore, different transport mechanisms and regimes can be obtained according to the nature of the electron-atom interaction (repulsive or attractive), the thermodynamic state of the gas either close to or removed from its critical point, and the amount of disorder inherent to the fluid [1].

Typically, at low density and high temperature, electrons are quasifree. Their wave function is pretty delocalized and the resulting mobility is large. They scatter elastically off the atoms of noble gases in a series of binary collisions and the scattering process is basically determined by the interaction potential through the electron-atom scattering cross section. The mobility can be predicted accurately by the classical kinetic theory [2].

At higher densities, and possibly, at lower temperatures, electrons may either remain quasifree with large mobility (as in the case of argon), or they can give origin to a new type of state that is spatially localized inside a dilation of the fluid. In this case the mobility is very low because the complex electron plus fluid dilation moves as an unique, massive entity. This, for instance, happens in He and Ne. The main difference between the two cases is that in the former the electron-atom interaction is attractive (Ar) and in the latter is repulsive (He and Ne) [3,4].

The simplest model to describe the behavior of electrons in a dense, disordered medium is the hard-sphere gas and a practical realization of this system is represented by He. In He the electron-atom interaction is pretty well described by a hard-core potential and the scattering cross section is fairly large and energy independent. It is well known that the charge transport proceeds via bubble formation in liquid He at low temperature [5–7]. In gaseous He at low temperature the mobility shows a drop of several orders of magnitude when the density is increased from low to medium values [8–12]. This drop has been intepreted in terms of a continuous transition from a transport regime where the excess electrons are quasifree to a region where they are localized. There is still controversy about the nature of the localized electron states in the gas, whether they are localized in bubbles, as in the case of the liquid, or whether they are localized in the Anderson sense [13]. In the latter case, the electron wave function decays exponentially with distance owing to multiple scattering effects induced by the disorder of the medium [14].

Owing to these considerations, it is interesting to investigate the localization transition at higher temperatures. Therefore, we have measured the mobility of excess electrons in dense He gas at temperatures $(26 < T < 77)$ K. By assuming that electrons are localized in dilations of the gas, a simple quantum-mechanical model provides a good semiquantitative description of the observed behavior of the mobility.

II. EXPERIMENTAL DETAILS

The mobility measurements have been carried out by using a swarm technique in a pulsed townsend photoinjection apparatus we have been exploiting for a long time for electron and ion mobility measurements [15–17]. A schematics of the apparatus is shown in Fig. 1. Briefly, a high-pressure cell (CN), that can withstand pressures up to 10 MPa, is mounted on the cold head of a cryocooler inside a triple-shield thermostat. The cell is operated between 25 and 330 K. Temperature is stabilized within 0.01 K.

A parallel-plate capacitor, consisting of an emitter (E) and a collector (C), is contained in the high-pressure cell and is energized by the high-voltage generator V. A digital voltmeter (DV) reads the voltage. The distance between the two plates delimits the drift space. An electron swarm is produced by irradiating the gold-coated quartz window placed in the emitter with the vacuum ultraviolet (VUV) light pulse of a Xe flash lamp (FL). The amount of emitted charge ranges between 4 and 400 fc, depending on the gas pressure and on

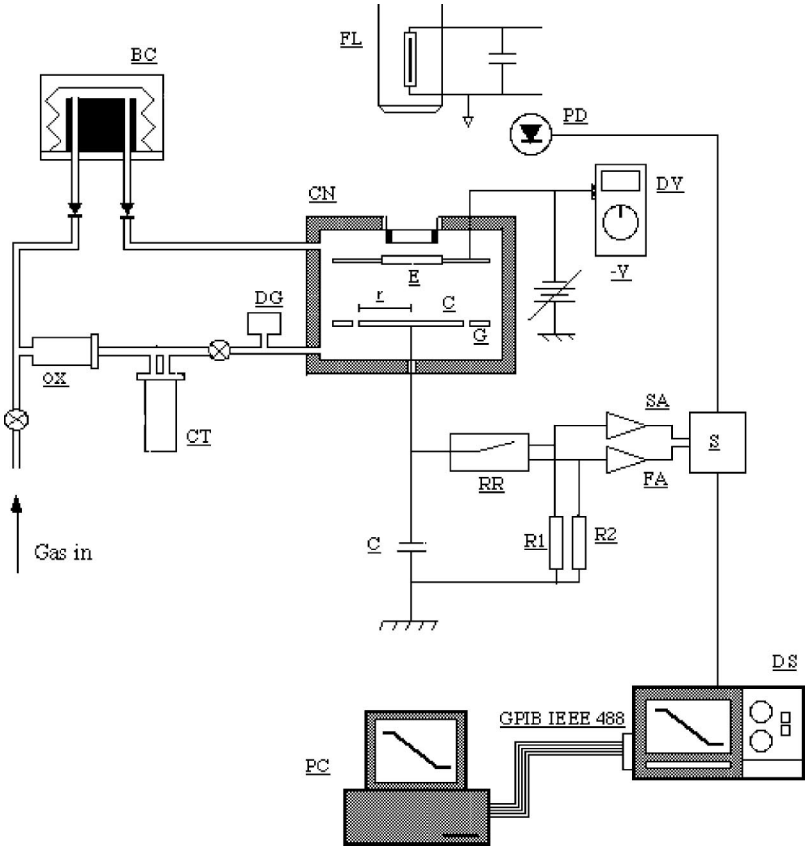

FIG. 1. Schematics of the experimental apparatus. See the text for a description.

the applied electrical field strength. Under its action, the charges drift towards the anode inducing a current in the external circuit. The current is integrated by the analog circuit RC in order to improve the signal-to-noise ratio. Two different operational amplifiers (SA and FA) are used depending on the duration of the signal. This is recorded by a high-speed digital transient analyzer (DS) and is fetched by a personal computer for the analysis of the wave form.

Ultrahigh purity He gas with an impurity content, essentially oxygen, of some ppm is used. The impurity content is reduced to a few ppb by circulating the gas in a recirculation loop driven by a homemade bellow circulator (BC) that forces the gas to flow through an Oxisorb cartridge (OX) and a LN2-cooled active-charcoal trap (CT).

The induced signal wave form of electrons drifting at constant speed is a straight line, and the drift time is easily determined by the analysis of the wave form. The overall accuracy of the mobility measurements is $|\Delta\mu/\mu|\approx 5\%$ [18].

III. EXPERIMENTAL RESULTS

In Fig. 2, we show the observed zero-field mobility μ_0 in He at $T\approx 26$ K. The present data are compared with literature data for $T=4.2$ K [8–10] and for $T=20.3$ K [11]. At $T=26$ K, μ_0 exhibits the same qualitative behavior observed earlier at much lower temperatures. As the gas density increases, μ_0 decreases by nearly five orders of magnitude. The continuous transition from the low-density, high-mobility region to the high-density, low-mobility one is interpreted as the progressive depletion of extended or delocalized states and the consequent formation of localized states [13,19]. These are assumed to consist of an electron trapped into a cavity in the fluid. This cavity is referred to as an electronic bubble [1,19].

A similar physical process has been observed also in liquid [20,21] and gaseous neon [16]. In gaseous neon, the μ_0 data resemble closely to those shown in Fig. 2 and the interpretation of the electron mobility behavior in neon, as due to electron localization in cavities, has been confirmed by quantum-mechanical molecular dynamics calculations [22].

The dynamics of the localization process, though not investigated experimentally, is quite clear [21,23,24]. However, even though the localization process were of the

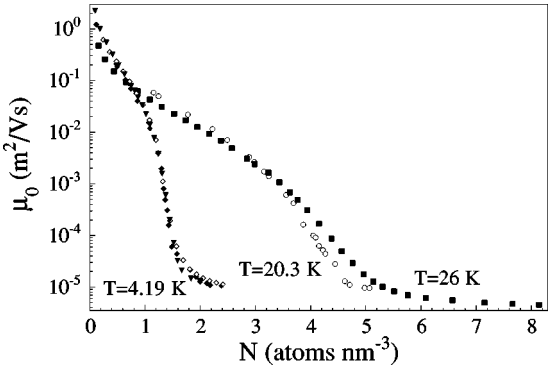

FIG. 2. Experimental zero-field mobility μ_0 as a function of the gas density. $T=26$ K: present work; $T<26$ K: literature data at $T\approx 4$ K [8–10] and at $T\approx 20$ K [11].

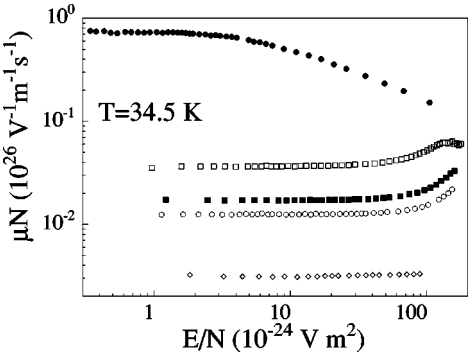

FIG. 3. Density-normalized mobility μN as a function of the reduced electric field E/N at $T=34.5$ K for several densities: $N=0.154$, 4.17, 4.66, 4.83, and 5.56 atoms nm^{-3} (from top).

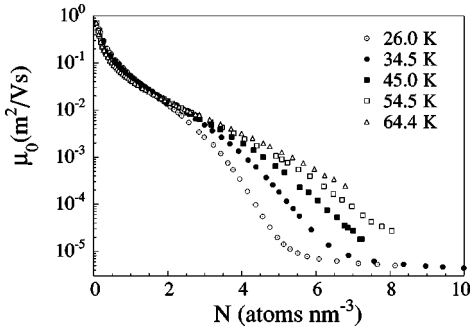

FIG. 4. Zero-field mobility μ_0 as a function of N for $T=26$, 34.5, 45.0, 54.5, and 64.4 K.

Anderson-type [13], i.e., self-localization of electrons with energy below the mobility edge as a consequence of self-interference of their wave function induced by the medium disorder, nonetheless electrons might wind up by forming electron bubbles because of the repulsive electron-medium interaction and medium compliance. The real existence of such electron bubbles has been also confirmed experimentally by infrared absorption spectra in liquid He [25,26].

Once all of the electron states are localized, the resulting μ_0 is not zero because the gas is compliant enough to allow the large complex structure made of an electron plus the associated bubble to diffuse slowly and drift under the action of an external electric field [1].

The main difference between the present data and those at lower temperatures is that the transition to low mobility states is shifted to larger density values. At $T=4.2$ K the transition can be considered complete at a density $N\approx 2$ atoms nm^{-3}. At $T\approx 20.3$ K the final state is reached for $N\approx 4.8$ atoms nm^{-3}, while at $T=26$ K in our experiment this density has moved to $N\approx 6.2$ atoms nm^{-3}. This is even more evident at higher temperatures.

It is clear that the formation of localized states is not related to the presence of a nearby critical point (the critical point of He is at $T_c\approx 5.2$ K and density $N_c\approx 10$ atoms nm^{-3}). It rather seems related to the competition between the thermal energy of electrons and the free energy of localization. Therefore, it appears reasonable that the localization transition shifts to larger densities for higher temperatures in order to achieve more favorable free energies.

The localization transition can be noticed also by observing the electric field dependence of the mobility. In Fig. 3, we plot the density-normalized mobility μN as a function of the reduced electric field E/N at $T=34.5$ K for several densities.

At small N and low E/N, electrons are in near thermal equilibrium with the gas atoms and μN is constant. As E/N increases, μN decreases, eventually reaching the $(E/N)^{-1/2}$ dependence expected on the basis of the classical kinetic theory because the scattering rate increases with the electron kinetic energy [2].

At high N, μN is very low and practically independent of E/N, at least for the highest electric fields of the present experiment (up to ≈ 7 kV/cm). At such densities, almost all of the electrons are localized in bubbles. Even the highest electric field reached in the experiment is not large enough to heat up such massive objects. The electronic bubbles, therefore, remain in equilibrium with the gas atoms.

At intermediate values of N, the behavior of μN is quite complicated. At small E/N, μN is constant, while at larger E/N, μN reaches a maximum and finally, at even larger E/N, it meets the classical $(E/N)^{-1/2}$ behavior. The same superlinear behavior of the drift velocity of electrons in dense He gas was observed also at very low temperatures but no interpretation was given then [10].

The complex density and field dependence of the mobility previously described for $T=34.5$ K is observed at all investigated temperatures.

This observed behavior can be easily interpreted in terms of the formation, at large N, of electron states that are self-trapped in partially filled bubbles. These are very massive and have low mobility. By increasing the electric field strength, bubbles may be either destroyed or their formation may be inhibited, so that electrons are again free and very mobile. The same behavior of μN as a function of E/N has been observed also in neon gas and the same interpretation of the data has proven successful [16]. Moreover, there is experimental evidence [21] that quasifree, highly mobile electrons do indeed exist at high electric fields even in liquid Ne, where they are usually localized in bubbles at small fields.

In Fig. 4, the zero-field value μ_0 of the mobility μ is shown as a function of the density N for the investigated temperatures. In this figure the shift of the localization transition to larger N for increasing T is clearly shown. For $T>45$ K the transition has not been tracked down completely because the pressure required to reach such large N values exceeds the capacity of our apparatus ($P\leqslant 10.0$ MPa). Nonetheless, it is evident that the localization phenomenon occurs also at high temperatures provided that the density is large enough.

IV. DISCUSSION

A description of the observed behavior of μ_0 as a function of N is very difficult. In fact, it must deal with the mobility of two charge carriers, the extended and the localized electron,

and it must also treat correctly the probability of occupation of the two states as a function of the density.

A further difficulty is that, although the mobility of the localized electron, i.e., of the bubble, is rather well described by the simple Stokes hydrodynamic formula, $\mu_0 = e/6\pi\eta R$, where η is the gas viscosity and R is the bubble radius [19,27], the description of the mobility of the extended electron states is still rather controversial, also because the localization transition is not as sharp as desired, as, for instance, in the case of Ne [16].

Several theoretical models for the description of the quasifree electron mobility in dense noble gases have been devised on the basis of the Boltzmann formalism of kinetic theory [3,13,14]. Their common feature is the realization that the multiple scattering effects concur to dress the electron-atom scattering cross section.

In particular, it has also been suggested [28] that, when the ratio between the electron thermal wavelength λ_T and its mean free path ℓ_c is $\lambda_T/\ell_c \approx 1$, the scattering rate diverges [13,28] and electrons get localized as a consequence of the interference of two scattering processes: the scattering off several different scattering centers and the time-reversed scattering sequence [29]. This model naturally introduces a mobility edge, an energy below which the electron wave function does not propagate.

Although this mobility-edge (ME) model describes quite well the electron mobility in dense He gas, it has two main drawbacks. The first one is that it works correctly only for He, because its scattering cross section is large and nearly energy independent. For Ne, for instance, it does not correctly describe the experimental data because of the strong energy dependence of the momentum transfer scattering cross section [30,31].

Moreover, it is well known that, in liquid He, electrons trapped in stable cavities within the fluid have been observed by infrared spectroscopy [1,25,26] and this observation has been confirmed also by quantum-mechanical molecular dynamics calculations [32], while the localized states, described in the ME model as those with energy below the mobility edge, are not only propagating but do not even reside in cavities. Even a static disorder produces localized electrons in this ME model. It is of course possible that after localization electrons could deform the fluid to produce bubble states, but the observed drop of mobility is not attributed here to bubble formation [13].

In view of these considerations, we adopt a simple model [33] that describes the formation of the self-trapped electron states as a process of localization in a quantum well. The mobility of the quasifree electrons is treated in terms of a different, heuristic model developed in our laboratory that encompasses the several multiple scattering effects present in the scattering process of an excess electron in a dense gas. We use such a model because it has given excellent agreement with the experimental data in Ne [15,16] as well as in Ar [34].

In addition to the usual thermal energy, electrons in the propagating state have a ground state energy $V_0(N)$ that depends on the density of the environment [35,36]. $V_0(N)$ consists of two contributions [37],

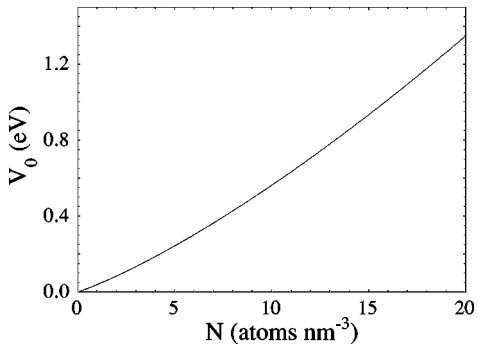

FIG. 5. Ground state energy $V_0(N)$ of a quasifree electron as a function of the gas density.

$$V_0(N) = E_K(N) + U_P(N). \tag{1}$$

U_P is a negative potential energy term arising from the screened polarization interaction of electrons with the gas atoms. $E_K(N)$ is a positive kinetic energy contribution due to excluded-volume quantum effects.

Owing to the small He polarizability, U_P can be neglected thus yielding $V_0 \approx E_K$. It has been shown [16,34,38] that E_K is quite accurately given by the Wigner-Seitz model

$$V_0 = \frac{\hbar^2 k_0^2}{2m}, \quad \tan[k_0(r_s - \tilde{a})] = k_0 r_s \tag{2}$$

as shown in Fig. 5. In Eq. (2), $r_s = (3/4\pi N)^{1/3}$ is the Wigner-Seitz radius, $\tilde{a}$ is the electron-atom scattering length, and k_0 is the ground state momentum of the electron. Owing to the fact that the electron-atom interaction is essentially repulsive, $V_0(N)$ is positive and increase monotonically with N. This means that the lower is the gas density, the lower is the ground state energy of a quasifree electron. This density dependence of V_0 is the main physical reason for electrons to favor regions of lower than average density.

V_0 fluctuates since thermally activated fluctuations of the density are present, and electrons can get temporarily localized in a virtual or resonant state above one such density fluctuation where the local density is lower than the average one [1,39].

If the electron-atom interaction is strongly repulsive (as in the case of He) and if the fluctuation is sufficiently deep, there can be formation of a self-trapped electron state, whose stability can be determined by minimizing its free energy with respect to the quasifree state.

We therefore assume that localized electron resides in a quantum square well of spherical symmetry. The well radius is R.

Since the gas has no surface tension and since the temperature is pretty high for He atoms to have significant thermal energy, we must allow for some He atoms penetrating into the cavity and dynamically interchanging with outside atoms. We thus assume the bubble to be partially filled with density $N_i < N$ and filling fraction $F = N_i/N < 1$. The electron is thus subjected to the following spherically symmetric potential:

$$\begin{cases} V_i & \text{for } r<R \\ V_0 & \text{for } r\geqslant R, \end{cases}$$

where V_i is defined as the ground state energy of an electron inside the bubble. Since $N_i<N$, $V_i<V_0$. The potential inside the bubble must take into account also the contribution of the polarization energy due to the outside gas. If the bubble were empty, the polarization energy $\mathcal{E}_P$ could be written as [33]

$$\mathcal{E}_P=-\frac{\alpha e^2}{2(4\pi\epsilon_0 R)}N. \tag{3}$$

Since the bubble is only partially empty, the polarization energy contribution can be written to first order as [16]

$$\mathcal{E}_P=-\frac{\alpha e^2}{2(4\pi\epsilon_0 R)}(1-F)N. \tag{4}$$

In this case the potential energy of the electron inside the bubble can be cast in the form

$$V_i=V_F+\mathcal{E}_P, \tag{5}$$

with $V_F=V_0(FN)$, i.e., the V_0 value at the density of the interior of the bubble.

A solution of the Schrödinger equation is sought for the lowest bound *s*-wave state, if it exists, of energy eigenvalue $\mathcal{E}_1$. Only the first eigenvalue is relevant because the temperature is quite low. If $\mathcal{R}(r)$ is the ground state solution of the radial Schrödinger equation, the function $f(r)=r\mathcal{R}(r)$ fulfills the radial equation

$$\left[\frac{d^2}{dr^2}+k_i^2\right]f(r)=0 \quad \text{for} \quad r<R,$$

$$\left[\frac{d^2}{dr^2}-k_o^2\right]f(r)=0 \quad \text{for} \quad r\geqslant R,$$

where $k_i^2=(2m/\hbar^2)(\mathcal{E}_1-V_i)$ and $k_o^2=(2m/\hbar^2)(V_0-\mathcal{E}_1)$.

By imposing the boundary conditions on the radial wave function at the bubble boundary for $r=R$, we obtain the eigenvalue equation

$$\tan X=-\frac{X}{(H^2-X^2)^{1/2}}, \tag{6}$$

with $X=k_iR$ and $H^2=(2m/\hbar^2)(V_0-V_i)R^2$. If X_1 is the solution of Eq. (6), then the energy $\mathcal{E}_1$ of the *s*-wave state is

$$\mathcal{E}_1=\frac{\hbar^2}{2mR^2}X_1^2+V_i. \tag{7}$$

The Schrödinger equation admits solutions if the well strength is such that $H^2\geqslant\pi^2/4$. This translates into a condition on a minimum bubble radius for the existence of a solution, namely,

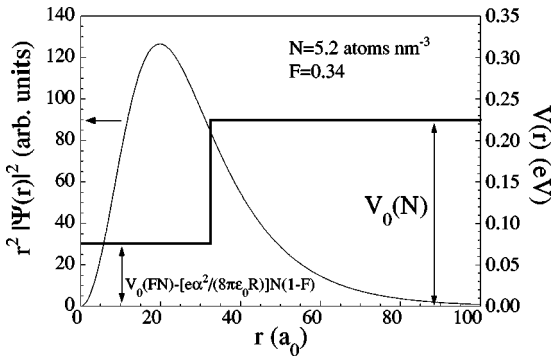

FIG. 6. Radial probability density for the *s*-wave ground state in the partially empty spherically symmetric square well. The distances are in units of the Bohr radius a_0.

$$R_0^2=\frac{\hbar^2\pi^2}{8m(V_0-V_i)}. \tag{8}$$

For each value $R>R_0$ the eigenvalue equation (6) is solved for X_1, and the eigenvalue $\mathcal{E}_1$ is calculated from Eq. (7) as a function of the gas density and of the filling fraction of the bubble.

In Fig. 6, we show the shape of a typical *s*-wave solution of the Schrödinger equation. The excess free energy of the localized state with respect to the delocalized one can be computed as

$$\Delta A=\mathcal{E}_1+V_i+W-V_0, \tag{9}$$

where W is the volume work, at constant T, required to expand the bubble and is given by [16]

$$W=\frac{4\pi}{3}R^3P\left[1-F-\frac{FN}{P}\int_{FN}^{N}\frac{P}{n^2}dn\right], \tag{10}$$

where P is the gas pressure.

In order to find the most probable state, ΔA is minimized with respect to the bubble radius and filling fraction. Rigorously speaking, the minimum excess free energy should be obtained by averaging ΔA over all atomic configurations leading to trapped electron states. This is a formidable task

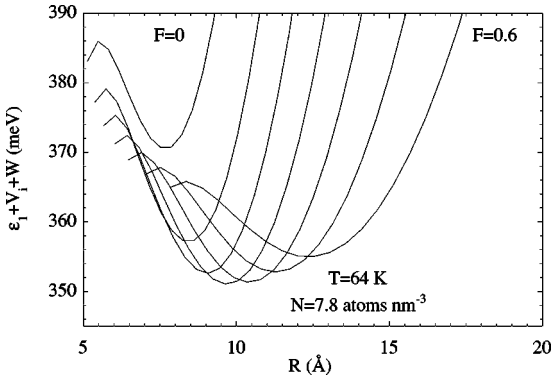

FIG. 7. Free energy of the localized state for given $T=64$ K and $N=7.8$ atoms nm^{-3} for filling fraction $F=0$, 0.1, 0.2, 0.3, 0.4, 0.5, and 0.6 as a function of the bubble radius.

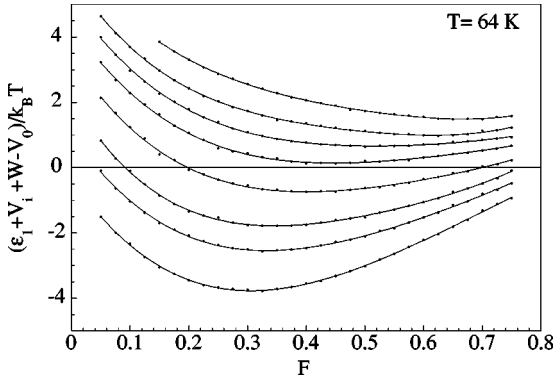

FIG. 8. The excess free energy, minimized with respect to the bubble radius, as a function of the filling fraction for $N=5$, 5.3, 5.5, 5.7, 6, 6.3, 6.5, and 6.8 atoms nm^{-3} for $T=64.4$ K (from top).

and therefore, to a first approximation, we adopt the optimum-atom-concentration fluctuation [19], i.e., that which causes the largest decrease of the system free energy as a consequence of electron trapping.

In Fig. 7, we show the free energy of the localized state $\mathcal{E}_1+V_i+W$ as a function of the bubble radius at fixed $T=64$ K and $N=7.8$ atoms nm^{-3} for several filling fraction values. The excess free energy values, minimized with respect to the bubble radius, are plotted in Fig. 8 as a function of the filling fraction for several N at fixed temperature.

For smaller densities, this excess free energy minimized with respect to bubble radius at constant N and T is a monotonically decreasing function of the filling fraction F. This means that the incipient bubble is not stable. It gets more and more filled until it disappears completely.

Stable states, marked by a minimum of the excess free energy, only appears at higher densities. Such stable states are now sought by carrying out a second minimization procedure of excess free energy as a function of the filling fraction F.

This double minimization procedure finally yields the optimum values of filling fraction F_B and bubble radius R_B, shown in Fig. 9 for $T=26$ K as a function of the gas density.

From Fig. 9 it can be seen that, at constant T, bubbles tend to become smaller and emptier as the density increases. The

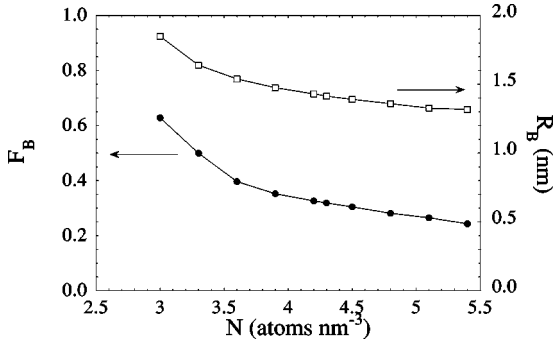

FIG. 9. Optimum equilibrium filling fraction F_B and radius R_B of the electron bubble for $T=26$ K as a function of N.

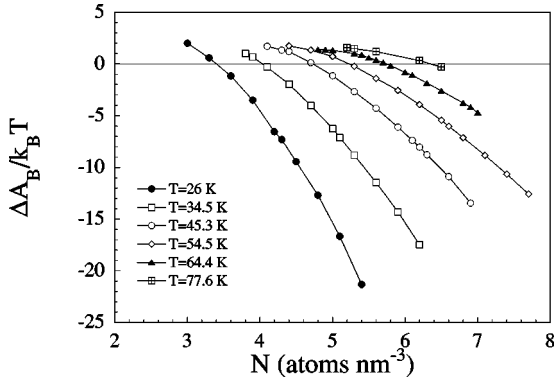

FIG. 10. Minimum excess free energy of the localized state as a function of N for several T.

optimum bubble radius $R_B \approx 1.0-1.8$ nm, depending on T and N, is compatible with the observed values in liquid He [1]. The density dependence of R_B and F_B at higher temperatures is similar. However, as a general trend, at high and constant densities, R_B decreases very slightly with increasing T while F_B increases quite a bit.

The values of the excess free energy corresponding to the optimum filling fraction and bubble, $\Delta A_B = \Delta A(R_B, F_B, N, T)$, are reported in Fig. 10. Bubble states start forming as soon as $\Delta A_B = 0$, but they are not stable against thermal fluctuations until $|\Delta A_B/k_BT| \gg 1$. For a given T, this condition is fulfilled only if N is large enough. Moreover, by inspecting Fig. 10, we see that a given value of ΔA_B is obtained at increasingly higher densities as the temperature is increased.

In Fig. 11, we show the values of density N^* where $\Delta A_B = 0$. At this density, localized and delocalized states are equiprobable. In agreement with the experimental observation on mobility, N^* increases with T. This means that bubbles become stable at larger N when T increases, both because electrons have more thermal energy and because the volume work to expand the bubble increases with the temperature.

Once the minimum excess free energy has been computed, the fraction of bubble and quasifree states is readily calculated as $n_B/n_F = \exp\{-\Delta A_B/k_BT\}$. The observed mobil-

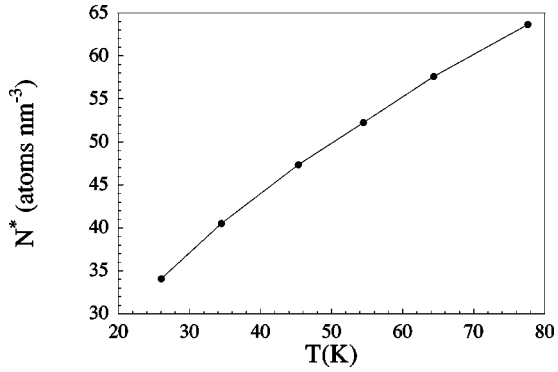

FIG. 11. Density values N^* where localized and delocalized states are equiprobable.

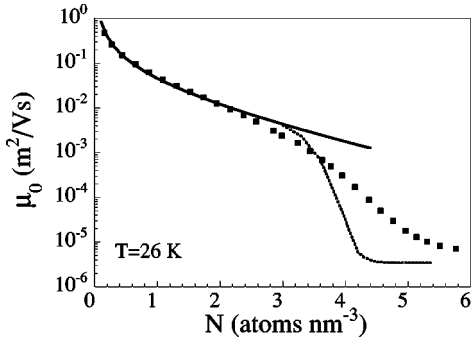

FIG. 12. Zero-field mobility μ_0 vs N for $T=26$ K. The solid line is the mobility of quasifree electron states. The dashed line is the weighted average mobility.

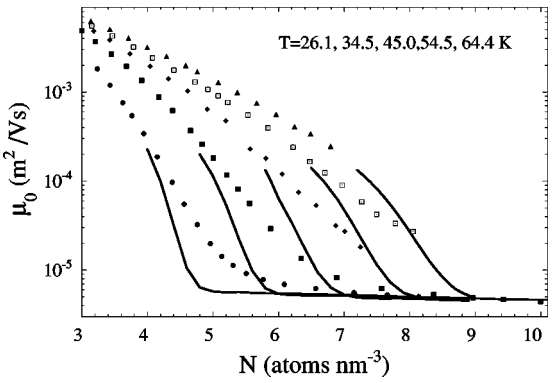

FIG. 13. Zero-field mobility μ_0 vs N for $T=26.1$, 34.5, 45.0, 54.5, and 64.4 K in the high-density region (from left). The solid line is the calculated average mobility.

ity is then a weighted sum of the contribution of the mobilities of the two states [19,40]. For the bubble state the semihydrodynamic mobility

$$\mu_B=\frac{e}{6\pi\eta R_B}\left[1+\frac{9\pi\eta}{4NR_B(2\pi mk_BT)^{1/2}}\right] \tag{11}$$

has been used [27], where η is the gas viscosity [41].

For the mobility of the quasifree states, we have used the results of the heuristic Padua model, succesfully exploited in Ne [15,16] and Ar [34]. The quasifree electron mobility can be written as [42]

$$\mu_F=\frac{4e}{3hS(0)}\lambda_T\lambda^*\exp(-\lambda_T/\sqrt{\pi}\lambda^*), \tag{12}$$

where h is the Planck's constant. $S(0)=Nk_BT\chi_T$ is the long-wavelength limit of the static structure factor and χ_T is the gas isothermal compressibility. $\lambda_T=h/\sqrt{2\pi mk_BT}$ is the thermal wavelength of the electron. Finally, λ^* is defined as

$$\lambda^*=\frac{1}{N}(k_BT)^2\int_0^\infty\frac{\epsilon}{\sigma_{mt}(\epsilon+E_k)}e^{(-\epsilon/k_BT)}d\epsilon, \tag{13}$$

where $\sigma_{mt}(\epsilon+E_k)$ is the momentum transfer scattering cross section evaluated at the electron energy shifted by the kinetic contribution E_k of the ground state energy shift V_0. We recall here that, for He, $E_k\approx V_0$. The exponential factor in Eq. 12 is due to O'Malley [14]. This model includes the three main effects of multiple scattering [34]: (1) the shift V_0 of the ground state energy of a quasifree electron in a medium of density N; (2) the correlation among scatterers taken into account by the static structure factor $S(0)$ [43]; (3) the increase of the scattering rate due to quantum self-interference of an electron multiply scattered in a time-reversed sequence by the same scattering centers [29] and described by the O'Malley factor in Eq. (12).

In Fig. 12, we show the results of the model for $T=26$ K. The quasifree mobility in the low-density side is well described by the heuristic model and also the density where the localization transition occurs is reproduced with satisfactory accuracy. Similar results are obtained for the higher temperatures.

In Fig. 13, we show the experimental mobility in the high-density region for $26<T<64$ K with the average mobility at high density calculated according to the present model. This figure clearly shows that the present model quite accurately predicts the shift of the localization transition to higher densities when the temperature is increased, although it does not fit the data with great accuracy.

V. CONCLUSIONS

The electron mobility in dense He gas shows two distinct regimes at low and high N. At low N the states of the excess electrons are extended, while at high N electrons are localized in bubbles. Both states are present at all N, but bubble states become stable, at fixed T, only if N exceeds a certain value N^*. The measured mobility is a weighted sum of the contribution of the two kind of electrons, quasifree and localized.

A simple model of electron localization in a quantum square well explains the observed fact that the localization transition shifts to higher N as T increases. It also semiquantitatively describes the observed mobility. The agreement of the model with the data, however, is far from satisfactory. More sophisticated models, namely, those based on the so-called self-consistent-field approximation [19,44], where the density profile of the bubble is self-consistently calculated along with the electron wave function, can be used but their results are not very different from the present ones.

Among possible reasons to explain the discrepancy of the present model with the experimental data, there could be the fact that the bubble model is a simple two-state model and neglects the possibility that bubbles have a distribution of radii and filling fractions. Moreover, even the description of mobility of the quasifree electrons is not yet completely satisfactory.

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