PART I

Electromigration I
BASIC CONCEPTS IN ELECTROMIGRATION

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ABSTRACT

A review of basic concepts in electromigration is presented and recent theoretical developments are discussed. The microscopic origin of the driving force for electromigration is elucidated and, as an example, the driving force exerted on an impurity near a grain boundary is calculated in a jellium model. A connection is made between the calculated electromigration driving force and 1/f-noise measurements. Results of the first computer simulations of electromigration at the microscopic level are presented. It is found that dynamical recoil effects can lead to a significant enhancement of electromigration for light interstitial impurities.

INTRODUCTION

When an electronic current flows through a conductor, a migration of atoms occurs. This phenomenon, which is known as electromigration, is a subject of technological interest primarily because of the deleterious effects of electromigration on VLSI operation. Electromigration is also an extremely interesting topic from a fundamental point of view in that it is an irreversible process of atom transport caused by dynamical coupling between electron and atom sub-systems. As such, it involves a number of subtle issues concerning electron transport in the vicinity of a diffusing atom, and how the electron transport gives rise to a driving force on that atom.

Our understanding of electromigration has progressively improved as a result of the efforts of a number of workers over the past several years. As a result of this activity, a number of basic concepts have emerged, and the field of electromigration has been placed on a much firmer theoretical foundation. Here we shall review these basic concepts and discuss some of the interesting new theoretical developments in the field. Earlier reviews of the subject can be found in Refs. 1-3, and a discussion of the role of electromigration in metallization failure can be found in Ref. 4.

PHENOMENOLOGICAL EQUATIONS

We are interested in the response of electrons and atoms to an applied electric field. The linear phenomenological equations of transport theory provide the connection between particle currents and applied fields. For simplicity, we assume only one atomic species, and consider the system to be at constant temperature. (A more general treatment can be found in Ref. 1.) The electron particle current is denoted by \( J_e \) and the atomic, or ionic, particle current is denoted by \( J_i \). [The terms “atom” and “ion” are used interchangeably; they both refer to the entity consisting of a nucleus plus the local electron cloud (core + conduction electrons) surrounding the nucleus.] The linear phenomenological equations have the form

\[
\begin{align*}
J_i &= -L_{ii} \nabla (\mu_i + q \Phi) - L_{ie} \nabla (\mu_e - e \Phi) \\
J_e &= -L_{ei} \nabla (\mu_i + q \Phi) - L_{ee} \nabla (\mu_e - e \Phi)
\end{align*}
\]

where the \( L \) coefficients are the generalized conductivities, \( \mu_i \) and \( \mu_e \) are the chemical potentials of the ions and electrons, respectively, and \( q \) is the intrinsic ionic charge. \( \Phi \) is the macroscopic electrical potential, which is related to the macroscopic electric field \( \vec{E} \).
according to \( \vec{E} = -\nabla \Phi \). The charge of the electron is \(-e\). The coupling between electron and ion currents is described by the \( L_{ie} \) and \( L_{ii} \) coefficients.

In the usual conductivity experiment there are no concentration gradients (\( \nabla \mu_i = \nabla \mu_e = 0 \)) and the electrons and ions undergo simple drift. Eqs. (1) and (2) reduce to Ohm’s law for each charged species. Eq. (1) can now be written as

\[
\vec{J}_i = L_{ii} q^* \vec{E},
\]

where we have defined the effective charge

\[
q^* = q - \frac{eL_{ie}}{L_{ii}}.
\]

Note that Eq. (3) is in the form of Ohm’s law for the ionic current, with the cross-coupling between electrons and ions contributing only to \( q^* \) through the \( L_{ie} \) term in Eq. (4). This allows us to say that the ions are driven in an electric field \( \vec{E} \) as though they have charge \( q^* \) rather than \( q \), with the difference \( q^* - q \) arising from dynamic electron-ion coupling.

Experimental results for \( q^* \) are traditionally expressed in terms of the effective valence \( Z^* \), which expresses \( q^* \) in units of the magnitude of the electron charge. By definition,

\[
q^* = Z^* e.
\]

Determination of this effective valence is the primary concern of experimental and theoretical studies in electromigration.

In the typical electromigration experiment, one deduces \( Z^* \) from a measurement of \( \vec{J}_i \) and \( L_{ii} \), and the subsequent use of Eq. (3). (The quantity \( L_{ii} \) is essentially the diffusion coefficient.) Experimentally, it is found that \( Z^* \) often differs dramatically from the nominal valence of the ion. Denoting the nominal, or “bare”, valence as \( Z \), and setting \( q = Ze \), we find that Eq. (4) implies that

\[
Z^* = Z - \frac{L_{ie}}{L_{ii}}.
\]

The fact that \( Z^* \) for metals is usually negative and large in magnitude indicates that the dynamic coupling described by the \(-L_{ie}/L_{ii}\) term is usually dominant.

**THEORY OF THE DRIVING FORCE FOR ELECTROMIGRATION**

The primary quantity of interest in electromigration theory is \( Z^* \), the effective valence of an ion for electromigration. In accordance with our previous discussion, \( Z^* \) is related to the driving force, \( \vec{F} \), exerted on an ion by

\[
\vec{F} = Z^* e \vec{E}.
\]

We now consider various theoretical approaches for the calculation of \( Z^* \). Along the way we develop the basic concepts underlying electromigration theory.

**Ballistic Model for \( Z^* \)**

In their pioneering work, Fiks\(^5\) and Huntington\(^6\) independently developed the ballistic model for the electromigration driving force. According to their model, there are two contributions to \( \vec{F} \). First, there is a direct contribution from the macroscopic electric field acting on the bare valence \( Z \). Accordingly, this so-called direct force can be written as \( \vec{F}_d = Z e \vec{E} \). Second, there is a contribution to \( \vec{F} \) due to the momentum transfer from the
electrons to the ion as the electrons scatter ballistically off the ion. This so-called electron-wind force, which we denote by $F_{\text{w}}$, is proportional to $-\sigma_{\text{tr}} J$, where $\sigma_{\text{tr}}$ is the electron-transport cross-section for scattering by the ion and $J$ is the electron current density. It is convenient to write this relation as $F_{\text{w}} = -K e J$, where $K$ is proportional to $\sigma_{\text{tr}}$. Then introducing the resistivity $\rho$ for the sample, and noting that $J = E/\rho$, one obtains

$$Z^* = Z - \frac{K}{\rho}. \quad (8)$$

For isolated impurity ions, one finds that $K = (n_0/N)\rho_1$, where $n_0$ is the electron number-density in the sample, $\rho_1$ is the resistivity due to the impurities and $N$ is their number density. Within this model, $Z^*$ can also be written as $Z^* = Z - n_0 \ell \sigma_{\text{tr}}$, where $\ell$ is the electron mean-free-path, which includes the effects of electron scattering by defects and phonons. From the latter form, it is easy to understand why for simple metals the wind force dominates the direct force. This is because $n_0 \ell \sigma_{\text{tr}}$, which is the number of electrons contained in a (fictitious) cylinder of length $\ell$ and cross-section $\sigma_{\text{tr}}$, is substantially larger than $Z$, which is the number of valence electrons contained in an atomic volume of a pure metal made up of atoms of the impurity species. (The cylinder volume is typically substantially larger than an atomic volume because $\sigma_{\text{tr}}$ is typically on the order of $1 \text{Å}^2$ and $\ell$ is on the order of 100 Å or more for temperatures of interest.)

The ballistic-model Eq. (8) is applicable for the free-electron model. In the case of general band-structure, Eq. (8) is still valid, but the quantity $K$ is no longer simply related to $\sigma_{\text{tr}}$. In further extensions of Eq. (8) to a parabolic hole-band, it has been suggested that the sign in front of $K/\rho$ be changed to account for the opposite sense of carrier motion. However, this is an over-simplification that is not likely to have more than qualitative validity. In fact, no satisfactory theoretical argument exists that supports a simple hole-wind expression for a localized hole in a metal. A further difficulty in generalizing the ballistic model is that for defect complexes involving more than one ion it is not obvious how the momentum transfer should be apportioned among the various ions in the defect complex, nor, for that matter, among the surrounding lattice atoms. In the case of vacancy diffusion, for example, it would not be meaningful to calculate a force exerted on a missing atom. Huntington and Grone realized this problem, and suggested that to estimate $Z^*$ for vacancy migration the quantity $K$ in Eq. (8) should be taken as one-half the value appropriate to an isolated ion. Despite these difficulties in generalizing the ballistic model, the model has the essential physics right: Electrons do transfer momentum to the ion and this is often the dominant effect in electromigration.

Polarization Charge Models

As an alternative to the ballistic model, Bosvieux and Friedel calculated the force on an ion from the polarization of the electrons in the applied field. They realized that the force exerted on the ion is given by the local electrostatic force acting on the bare ion. This result, which for closed systems is known as the Feynman-Hellmann theorem, can be derived from the definition of force as the time-rate of change of the momentum of the atom. In this picture, if $\delta n(\vec{r})$ is the part of the electron density that depends on the electric field and current, then that part of the force on the ion which depends on the electric field and current is given by

$$\vec{F} = Ze \vec{E} - \int \delta n(\vec{r}) \frac{\partial V_e}{\partial \vec{r}} \, d^3 r, \quad (9)$$

$$e$$
where \( V_0 \) is the potential energy of interaction between an electron and the bare ion, and \( \vec{R} \) is the position of the ion. The integral is over all space. For a point-charge model of an atom, \( V_0 = -\frac{Ze^2}{|\vec{r} - \vec{R}|} \).

Eq. (9) is exact, but to use it we need to find the perturbed electron density \( \delta n(\vec{r}) \). Bosvieux and Friedel separated \( \delta n(\vec{r}) \) into two parts, namely an electrostatic screening part and a dynamic screening part due to the electron current. The former part corresponds to the electron density involved in the screening of the direct force and the latter part corresponds to the electron density induced by the electron wind. They gave an argument that the electrostatic screening on an isolated ion is complete, i.e., the direct force vanishes. This part of their analysis has met with considerable skepticism, and we have shown that their analysis of the direct force cannot be correct. Specifically, if one considers the formal limit of very small \( Z \), then \( \delta n \propto Z \), and so by Eq. (9) \( \vec{F} \) equals \( 2Ze \overline{E} \) plus corrections of order \( Z^2 \overline{E} \). Indeed, as \( Z \to 0 \), the ratio \( \vec{F}/Ze \), by the very definition of the electric field in electromagnetic theory, must equal the pre-existing microscopic electric field at the location of the impurity ion. That is, in the \( Z \to 0 \) limit, the ion is simply playing the role of the “test charge” that operationally defines the electric field. This is a very fundamental point that is missed by those who still erroneously claim that the direct force vanishes. To deny the existence of the direct force \( 2Ze \overline{E} \) in the small-\( Z \) limit is to deny the existence of a microscopic electric field at the location of the ion before the ion is placed there.

Despite this inconsistency in the Bosvieux-Friedel treatment of the direct force, their calculation of the wind-force is correct, and represents a major contribution in the evolution of electromigration theory. In their approach, the wind-force arises from the total electron density due to the self-consistently screened pile-up of scattered electrons. This electron density, \( \delta n(\vec{r}) \), is related to the scattering-state wavefunctions \( \psi_k(\vec{r}) \) for electrons incident upon the defect. Explicitly,

\[
n(\vec{r}) = \sum_k f_k |\psi_k(\vec{r})|^2, \tag{10}
\]

where \( f_k \) is the electron distribution in the transport problem, and the sum is over all electron states \( k \). In the free-electron model (or “jellium model”) used by Bosvieux and Friedel, \( f_k \) is simply a shifted Fermi sphere in \( k \)-space.

The \( \psi_k(\vec{r}) \) are determined by a self-consistent calculation in which the potential energy of the electrons consists of the bare potential energy \( V_0 \) due to the ion and the electrostatic potential energy arising from \( n(\vec{r}) \). Using first-order perturbation theory, Bosvieux and Friedel obtained a wind-force expression which we can cast into the following form:

\[
\vec{F}_w = -\frac{em\tau \overline{E}}{12\pi^2\hbar^3} \int_0^{2k_F} |V(q)|^2 q^3 dq, \tag{11}
\]

where \( m \) is the electron mass, \( \tau \) is the electron relaxation-time, and \( V(q) \) is the Fourier transform of the statically screened electron-ion interaction in the absence of \( \overline{E} \). For the point-ion impurity, \( V(q) = -4\pi Ze^2/q^2 \kappa(q) \), where \( \kappa(q) \) is the dielectric function.

It is easily verified that expression (11) is identical to the wind-force expression obtained in the ballistic model for an isolated ion. The prime virtue of the Bosvieux-Friedel approach, however, is that it can be generalized to systems more complicated than an isolated ion in a free electron gas. Bosvieux and Friedel provide examples of models that might be used to describe various diffusion complexes. However, their ad hoc modelling and their use of perturbation theory for a weak point-ion impurity requires some improvement. This was provided by our pseudopotential-based theory.\(^9\)

The pseudopotential-based calculation of \( \vec{F}_w \) is similar to the Bosvieux-Friedel calculation except that it requires self-consistent solution of the pseudopotential form of the
Schroedinger equation. For the case of an interstitial ion, one again obtains the wind-force expression (11) but with \( V(q) \) replaced by the pseudopotential form-factor. For other diffusion complexes, such as a vacancy complex, one has rather simple expressions which involve the instantaneous positions of all the atoms in the complex. The calculated \( Z^* \) values for vacancy diffusion in simple metals were found to be in reasonable agreement with experiment, with discrepancies on the order of 25% being typical.

A number of rather sophisticated calculations of the wind force have appeared more recently. In virtually all of these calculations, the starting point is the one-electron version of Eq. (9) in which the bare electron-ion potential energy \( V_o \) is replaced by the screened electron-ion potential energy \( V_s \). That is,

\[
\vec{F}_w = -\int \delta n(r) \frac{\partial V}{\partial R} \hat{e}_r. \tag{12}
\]

Further, in this one-electron prescription,

\[
\delta n(r) = \sum_k g_k |\psi_k(r)|^2, \tag{13}
\]

where \( g_k \) is the off-equilibrium (\( \vec{E} \)-dependent) part of \( f_k \), and \( \psi_k \) are now the scattering states for electrons in the absence of \( \vec{E} \) (and the associated non-equilibrium charge pile-up). In this scheme, one can evaluate \( \psi_k \) directly from atomic scattering potentials for the equilibrium electron system; non-equilibrium electron aspects enter only in \( g_k \) and not in any self-consistent screening potential. An impressive state-of-the-art calculation based on this scheme has recently been performed by van Ek for hydrogen electromigration in transition metals. van Ek’s calculations make use of a KKR distorted-lattice approach due to Lodder. A justification for Eqs. (12) and (13) has been provided by analyses based upon the linear-response theory to be discussed later.

Local Transport Field and Landauer’s Analysis

In electromigration, the bare ion feels a driving force given by the microscopic electric field acting on the bare ion. This is consistent with Eq. (9) and with our earlier discussion. Sources of the local field in the vicinity of a microscopic scatterer during electronic conduction were first elucidated by Landauer in his classic 1957 paper. According to Landauer, a localized dipole is set-up around the scatterer as incident electrons pile-up there. This dipole, which he called the residual resistivity dipole (RRD), provides the necessary local electric field to enable the electrons to traverse the scattering region. In addition, the long-range fields set up by the RRDs centered at each defect provide the extra macroscopic field associated with the residual resistivity. The RRD thus provides a direct link between microscopic and macroscopic levels.

Landauer and Woo later applied the RRD analysis to electromigration. They found that the force exerted on the ion by the RRD is in the direction opposite to the electron wind force. They estimated that the RRD contribution, though formally higher-order in \( Z \), could be comparable in magnitude to the Bosvieux-Friedel wind force. The situation was clarified by Schaich who showed that expression (13) of the one-electron polarization-charge model contains the RRD provided that \( \psi_k \) is calculated to second-order in the scattering potential. Although Schaich considered a non-interacting gas, electron screening does not modify this conclusion, at least within the simplest screening approximations. According to Schaich and Sham, the RRD is automatically contained in the wind expression (11) if one replaces \( V(q) \) by the T-matrix. The situation
is not quite so simple, however, when the defect complex contains more than one scatterer or when interaction between RRDs is considered.

Landauer identified another source of the local field, which he called the carrier density modulation (CDM) effect. This effect arises from the changed carrier density in the vicinity of a defect ion. The electron density is locally increased for a positively charged bare impurity ion \(Z > 0\) as electrons are locally drawn in to neutralize the impurity. The additional carriers make it easier to overcome the background lattice scattering and so cause the local electric field to decrease. Analysis of the CDM effect led Landauer to a direct force given approximately by

\[
\tilde{F}_d = Z e \tilde{E} + \frac{1}{3} \frac{\Delta n}{n_0},
\]

where \(\Delta n\) is the extra electron density in the region of the ion, i.e., \(\Delta n = Z/\Omega_i\), where \(\Omega_i\) is the local volume of the ion. Note that unlike the Bosvieux-Friedel direct force, Eq. (14) correctly reduces to \(ZeE\) in the formal limit of small \(Z\). In general, the term \(\frac{1}{3} \frac{\Delta n}{n_0}\) implies that the direct force can be appreciably smaller than \(ZeE\) for positive \(Z\) (and \(\Delta n\)).

The local transport field was also calculated by Das and Peierls22, who performed a semi-classical analysis assuming the weak-charge limit (small \(Z\)). Because their analysis is based upon the Boltzmann equation, it does not adequately address the quantum mechanics of the electron response. However, their calculation does have the virtue of being entirely self-consistent. As such, it represents the first such calculation in electromigration theory. Their final result for the force on a point-ion in a free electron gas can be written as

\[
\tilde{F} = Ze\tilde{E} + \tilde{F}_w,
\]

where \(\tilde{F}_w\) is given by Eq. (11) except that in Eq. (11) \(2k_F\) in the upper-limit of the integral is replaced by a cut-off wavevector, which is taken to be on the order of the Thomas-Fermi wavevector. (The latter is given by \((e/h)\sqrt{4mk_F/\pi}\), which does turn out to be on the order of \(2k_F\) for alkali metals.)

The Das-Peierls analysis was later generalized to the quantum-mechanical case by Sorbello and Dasgupta23 who employed a density-matrix analysis. This analysis was subsequently improved upon by Schach24, and later applied to electromigration in a semiconductor inversion layer.25 In the limit of large mean-free path, the wind-force was found to approach the wind-force of the ballistic model. Again, the direct force approached \(Ze\tilde{E}\) in the formal limit of small-\(Z\), with corrections of order \(Z^2\tilde{E}\) also present.25-27

### Linear Response Theory

Kumar and Sorbello27 introduced the Kubo linear-response formalism into electromigration. In so doing, they cast the electromigration problem into a rigorous framework which allowed many-body techniques to be used. The formalism was later used by others20,24,28-36 and ultimately led to justification of the wind-force expressions (12) and (13). As far as the direct force is concerned, we were able to use this formalism to derive a tractable expression for a single impurity in a free electron gas.9 This direct-force expression involves simple Green functions, and was applied to a model for hydrogen impurities in metals. The results9 showed that the direct-force valence \(Z_d\) is neither zero (the Bosvieux-Friedel value) nor \(Z\) (the Huntington-Fiks value), but instead was typically on the order of 10%-30% smaller than \(Z\). The non-integral aspect of \(Z_d\) is borne out by experiment.9

In trying to understand the physical origin of the direct force, it is important to realize that the direct force is not due to a simple equilibrium polarization response of the neutralizing electron cloud surrounding an ion. Indeed, if that were the case, the direct force would vanish.9 Rather, it is worth keeping Landauer’s CDM picture21 in mind, and
realize that the direct force involves electron currents that are trying to push through the region of the ion in the presence of background scattering. If we imagine an extremely small, homogeneous system having dimensions much smaller than the background mean-free-path, there would in fact be no direct force. In such “mesoscopic” systems, the local field arises only from coherent scattering by the ion, or ions, in the system. Since these scattering events are fully included in the wind-force, there would be no additional direct-force contribution to consider in such systems.  

Z* in Grain Boundaries

As an interesting application, we consider an impurity near a grain-boundary in a simple metal. As usual, the driving force is given by the sum of the direct force and the wind-force. For a simple metal, the wind-force dominates, and we therefore restrict our attention to this quantity, which we can calculate using Eqs. (12) and (13). In the simple jellium model which we employ, we regard the grain-boundary to be a disordered region confined to the slab-like domain $0 < z < d$, where $d$ is the grain-boundary thickness. Since on average there are fewer atoms in this region than in the regular lattice, the grain boundary region presents a repulsive potential to the electrons. We take this potential energy to be a constant, $U$, for $0 < z < d$ and to vanish elsewhere. For the impurity, we use a delta-function potential energy $V(\mathbf{r}) = \delta(\mathbf{r} - \mathbf{R})$.

Assuming the incident electron current to be parallel to the grain-boundary, we can calculate $\psi(\mathbf{r})$ to second-order in $v$ and to all orders in $U$. The wind-force can then be found from Eqs. (12) and (13). The result is shown in Fig. 1 for $U/E_F = 0.1, 0.5$ and 0.99, where $E_F$ is the Fermi energy. The displayed wind-force is normalized so that it equals unity far away from the grain boundary, i.e., where the wind-force approaches its value for an isolated impurity in bulk metal. In our calculations we have chosen $d = 4$ a.u. (1 a.u. = 0.53 Å).

Figure 1. Wind force on an impurity near a grain boundary.
The results in Fig. 1 indicate that the wind-force is strongly diminished when the impurity is inside the grain boundary. Even for a relatively weak grain-boundary potential, the wind force inside the grain boundary is on the order of 20% smaller than the value just outside the grain-boundary. We have also calculated the normalized, local electronic current in the absence of the impurity; the current curves are found to be very similar to the wind-force curves.

Our calculations suggest a connection between electromigration and $1/f$-noise. It is easy to show that in the model being considered, the wind-force is proportional to the scattering cross-section presented by the impurity. When the impurity jumps from the region outside to the region inside the grain boundary, there is a cross-section change $\delta \sigma_{tr}$. Such jumps give resistivity fluctuations which can be observed as $1/f$-noise. Recent $1/f$-noise experiments can be explained provided that $\delta \sigma_{tr}/\sigma_{tr} \sim 0.20$. This value is in fact consistent with our calculated wind-force for an impurity jumping in and out of a grain boundary.

**Dynamical Effects in Electromigration**

In the standard picture which we have been assuming up to now, we have considered the electromigration driving force as if it were an adiabatic quantity. That is, we have neglected atomic recoil and have calculated $\vec{F}$ as if the impurity were infinitely massive. In this standard picture, the force so determined is then to be used as a static perturbation which tilts the atomic potential and causes an atomic current. In reality, the electron-wind force is a series of impulses due to individual electron collisions, and the atom will generally recoil and absorb energy in these collisions. For classical diffusion mechanisms, it is the energy transfer that is responsible for electromigration. But energy transfer depends on the velocity of the atom at the time of collision, the transfer being larger when the atom is moving at a larger velocity. This in turn leads to an atomic current that is enhanced over the value calculated in the purely adiabatic picture. This enhancement is expected to be a significant effect for the case of light interstitial impurities such as hydrogen.

To test this theory we have performed the first numerical simulations of the electromigration process. We consider the classical diffusion regime, and take our atom to be located in a harmonic oscillator well. The potential is $U(x) = \frac{1}{2} M \omega_0^2 x^2$ where $M$ is the atom mass and $\omega_0$ is the angular frequency of oscillation. A jump is taken to occur when $|x|$ is greater than some critical displacement $x_c$, which we take to be 1 Å. The atom is subjected to thermal collisions at an average frequency $\nu_{th}$, and these collisions randomize the atom velocity towards the equilibrium Boltzmann distribution. In addition, the atom is subjected to collisions with the electrons, at an average frequency $\nu_{el}$, which is proportional to the product of electron current and the scattering cross-section. (We estimate $\nu_{el} \sim 10^9 s^{-1}$ when $J \sim 10^6 A/cm^2$ and $\sigma_{tr} \sim 1 A^2$.) All collisions are taken to occur randomly via a Poisson process, i.e., there is an exponential distribution of waiting times for thermal collisions and for electron collisions. For simplicity, we assume a constant momentum transfer $p_e$ for each electron collision.

Calculations have been performed for a hydrogen atom and also for an aluminum atom. To make the computations manageable, we have chosen a rather small activation energy $E^* (E^* = \frac{1}{2} M \omega_0^2 x_c^2 = 0.1 eV)$, and have chosen a temperature of 300 K. The results for the net atomic current are shown in Fig. 2 for the case that $\nu_{th} = 10.0$ and $\nu_{el} = 0.1$, where $t_p = 2 \pi / \omega_0$ is the oscillation period. The net atomic flux (the number of jumps/sec in the direction of the electron wind) is shown as a function of the momentum transfer per collision. (The flux is in arbitrary units, chosen differently for H and Al so that the net...