ENHANCED EFFECTS OF STARLIGHT ON THE INTERSTELLAR MEDIUM

Humberto Gerola and Robert A. Schwartz
New York University

Received 1975 August 6

Revised

(NASA-CR-145707) ENHANCED EFFECTS OF STARLIGHT ON THE INTERSTELLAR MEDIUM (New York Univ.) 20 p HC $3.50 CSCL 03B

N76-12919

Unclas
G3/90 01977

ABSTRACT

The photodesorbtion of molecules and atoms from the surfaces of interstellar grains can be an important source of heating for the interstellar medium and the origin of instabilities which may separate grains and gas. For low densities, the force exerted on the grains is proportional to the gas density and independent of the radiation intensity; for high densities, it is proportional to the radiative flux and independent of the gas density. This force may act differently on grains of different sizes. The photoelectric effect may also be an efficient mechanism for the separation of gas and dust in diffuse clouds.

I. Introduction.

The non-ionizing radiation from stars has an energy density comparable to that of the other components of the interstellar medium, but its effects on the overall structure of HI regions have generally been thought to be small. We point out that this may not be the case, and that starlight, through its interaction with atoms or molecules adsorbed on the surfaces of grains, may play an important role in many parts of the interstellar medium. In this note, we present simple physical arguments which suggest the circumstances under which the effects of starlight may be important, and which will enable us to estimate the contribution of this process to the heating and the stability of the interstellar medium.

In Section II we discuss how the photodesorption of atoms or molecules adsorbed on the surface of interstellar grains may be a mechanism for the strong coupling of the grains to the stellar radiation field. In Section III we investigate the role which this process plays in the separation of gas and dust in the interstellar medium, and show that the canonical assumption of a constant ratio of gas to dust, while it may be true in some average sense, cannot be true for individual clouds. We shall exhibit an instability which occurs in uniform gas-dust clouds which spontaneously separates the two, even in the presence of a uniform stellar background. In Section IV we show that this process may also be an important heating mechanism for diffuse interstellar clouds which is not affected by the attenuation of radiation in the cloud. Finally, in Section V the different consequences of these results are briefly discussed.

II. Photodesorption and the Force on Single Grains.

Our discussion begins with the recognition by Reddish (1971) that, if near UV photons can liberate adsorbed molecules from the surfaces of interstellar grains, the momentum transferred to the grains will be enhanced by a large factor over that due to radiation pressure alone. The reader may observe a similar effect in a device usually known as the Crookes radiometer, which revolves in the opposite sense from that which one would expect from radiation pressure. (A Crookes radiometer in an extremely high vacuum goes around the "right" way.)

Although the photodesorption of atoms or molecules from solid surfaces has been known experimentally for a long time, it still is not clear which are the basic mechanisms responsible for it. Watson and Salpeter (1972) argue that the adsorption of a photon involves the transition to an excited electronic state of the adsorbed atom or molecule. By the Franck-Condon principle the transition occurs at a constant value of the internuclear distance (i.e. the distance from the surface to the adsorbed particle). Since the adsorption potential for the excited state is likely to be quite different from that of the ground state, the difference will usually be greater than the adsorption energy. If the particle proceeds to a distance of the order of several Angstroms before returning to the electronic ground state, it will leave the surface.

For sites of enhanced binding, on which particles are chemisorbed, however, the electron may be excited into a state which is equivalent to one of the electronic states in the solid. If the overlap between this state and the states inside the solid is large, the probability of desorption ε will be small; if the excited state is localized, ε may be large even for chemisorbed species. For highly dissimilar species bound to the crystal, the overlap is smaller, because of the difference in the adsorption bond length and the lattice constant.

The argument given above also leads us to assume that most of the photon energy is given to the desorbed particle, since the energy transfer to the lattice is expected to be small. If the desorbed particle leaves the surface in an excited electronic state, this state will be of a well-defined energy, and the difference between the photon energy and that of the desorbed excited state will also be precisely defined, so the difference between the photon energy and that of the excited state will appear in the only available form, as kinetic energy of the particle. We are therefore led to assume that a large fraction of the photon energy appears as kinetic energy of the desorbed particle. Other mechanisms that have been proposed for photodesorption are listed by Greenberg (1973). Recently, Gersten, Janow and Tzoar (1975) have made a detailed theoretical study of the photodesorption of single atoms. They estimate a quantum yield for photodesorption from a monolayer on a surface of $\approx 2 \times 10^{-3}$ in agreement with some earlier experimental work by Fabel, Cox and Lichtman (1973). Of particular interest here is the possibility, mentioned by Greenberg, that the molecule goes up to an excited electronic state, returns to the ground electronic state, but in an excited vibration-rotation level. This internal kinetic energy may be converted to translational energy away from the surface by collision with surface atoms. In this case the energy of the ejected molecule is related to the vibrational levels of the molecule rather than to the energy of the photon. Present experiments do not allow to discriminate among the different possibilities; Greenberg, however, from experimental data on several molecular species was able to estimate that the relative quantum efficiency is very near one, in accord with the arguments of Watson and Salpeter, so that an appreciable fraction of photons absorbed lead to desorption of molecules. The force exerted on a grain depends on the square root of the kinetic energy of the molecule; therefore, at this stage, we assume that the kinetic energy of the ejected molecule is essentially that of the photon.

A grain in which a fraction f of the surface area σ is covered by adsorbed atoms or molecules of mass m, experience a force

$$F_g = B\sigma f m^{\frac{1}{2}} \varepsilon \int_0^\infty H_\nu (h\nu)^{\frac{1}{2}} d\nu , \qquad (1)$$

where H_ν is the flux of starlight energy and ε the probability that a photon will release a particle from a fully covered surface. This estimate is under the assumption that the particles desorbed leave the grain isotropically; it only changes slightly for other angular distributions, B is an efficiency factor, near unity, which includes the effect of diffraction. If we take $f = 1$, $\varepsilon = 1$, and m to be the mass of a hydrogen molecule, this force is about 3×10^4 times stronger than the radiation pressure. It is therefore apparent that this grain jet effect, as we may call it, may be important if ε is large even if only a small fraction of the grain is covered by adsorbed atoms or molecules, or, vice versa, with a smaller ε on a completely covered grain.

The photodesorbed particle and the grain carry equal and opposite momentum; the net momentum transferred to the gas by this process is thus zero, if we neglect the momentum of the initial photon. The grain jet effect thus causes motion of the grains <u>through</u> the gas; the radiation exerts a force much larger than the radiation pressure on a small fraction of the mass of the gas-dust system. This means that the instabilities we shall find as a result of this force will grow very fast compared to the dynamical timescale associated with the gas itself.

In order to estimate f, we must consider the processes by which adsorbed particles stick to and leave the surface of the grains. An extensive discussion of these processes has been given by Hollenbach and Salpeter (1971, hereafter referred to as HS). They considered the process by which molecules stick to the grain, and the rate of thermal evaporation, but not photodesorption.

The average time for a particle to evaporate from the surface of a grain is $t_{ev} = (1/\nu_o) \exp(+D/T_g)$ in the notation of HS, where D is the adsorption energy, T_g the grain temperature and ν_o the characteristic vibration frequency of the adsorbed particle. The grain temperatures and the adsorption energies, of course, vary greatly according to the nature of the grains. For example, the adsorption energy of H atoms on a perfect ice surface was estimated by HS to be about 450° K, and for H_2 about 550° K. They estimated similar adsorption energies for graphite surfaces. However, binding energies of both atoms and molecules to defect sites on the grain surfaces can be many times this energy, up to chemisorption energies of several electron volts on highly reactive sites. If particles are adsorbed, or chemisorbed, on such enhanced sites, they will remain so for all reasonable grain temperatures, thus setting a lower limit to the value of f. Unfortunately, we do not know how to estimate the number of enhanced sites which may be present on a grain; we therefore shall not include the contribution of such sites in our estimates of f. It must be remembered that, if there is an appreciable number of such enhanced sites on a typical grain, the occupation number, f, will be larger and the grain jet effect correspondingly stronger.

The rate at which molecules stick to the surface of the grain is $t_s^{-1} = (\sigma \tilde{v} n S)$, where n is the gas density and S is a sticking probability, estimated by HS to be about 0.5. $\tilde{v}$ is a somewhat complicated function of the gas-dust velocity, but has the simple limiting forms $\tilde{v} = c_s$, the isothermal sound speed, if the dust is nearly stationary with respect to the gas, and $\tilde{v} = \frac{1}{2} v_{gas-dust}$ for supersonic grain velocities. Th. rate at which grains evaporate is $f\, t_{ev}^{-1}$ and the rate at which they are photodesorbed is

$$t_{\gamma}^{-1} = f\, \varepsilon \cdot 4\pi\sigma \int_0^{\infty} \frac{J_\nu}{h\nu} d\nu \equiv f\, \varepsilon\, \sigma\, \phi\ . \qquad (2)$$

We may determine f by equating the rate of sticking to the sum of the evaporation and photodesorption rates. As was shown by HS, the evaporation rate becomes large if T_g exceeds

$$T_c = D/\ln(t_s \nu_o) \sim 15^\circ\ K \quad . \qquad (3)$$

In that case, f will be negligibly small, and we may neglect the process under discussion. Estimates of free-space grain temperatures (e.g. Leung, 1975) give temperatures for silicate grains of 13° K, for graphite grains of 30° K, for silicate-ice grains of 13° K and for silicone carbide grains of 48° K. It is thus apparent that, for at least some types of grains which have been proposed, we may neglect evaporation under typical interstellar conditions.

If we neglect evaporation from the grains, then we may estimate f as simply

$$f = (S/\varepsilon)\, n\, \tilde{v}/\phi\ . \qquad (4)$$

This form is, of course, applicable only if $f \ll 1$. If f is greater than one by this prescription, we shall set it equal to one for purposes of our estimates. f is nearly one for densities greater than n_c, the critical density,

$$n_c = (\phi/\tilde{v})\,(\epsilon/S) \quad . \tag{5}$$

This density separates the regimes in which the grain jet effect has qualitatively different dependence on gas density and on the radiation field.

If $n < n_c$, the force on each grain is given by

$$F_g = n\, v\, \sigma\, m^{\frac{1}{2}}\, S < (h\nu)^{\frac{1}{2}} > \cdot\, \delta \ , \tag{6}$$

where δ, the anisotropy of the radiation field, is defined as

$$\delta = \frac{1}{4\pi\sqrt{2}} \cdot \frac{\int_0^\infty H_\nu/(h\nu)^{\frac{1}{2}}\, d\nu}{\int_0^\infty J_\nu/(h\nu)^{\frac{1}{2}}\, d\nu} \ . \tag{7}$$

This force has several remarkable properties. It is independent of the intensity of the radiation field, and depends only on its degree of anisotropy. This behavior is due to the assumption that evaporation can be neglected; in that case, f is inversely proportional to the flux of photons, so the number of photodesorptions per second is constant (and equal to the number of atoms sticking to the grain). For the same reason, the force is independent of ϵ, and proportional to n.

The drag force on a grain (moving through neutral gas) is also proportional to n, and to the square of the gas-grain velocity for supersonic motion. Thus, the velocity of the grains with respect to the gas is approximately $1/2\pi$ times the velocity with which the particles are ejected from the grain; for typical interstellar photon energies of 6 e.v. this corresponds to $\tilde{v} = 8$ km/sec. $\cdot\, \delta$.

In the opposite limiting case, $n \gg n_c$, the force is

$$F_g = m^{\frac{1}{2}} \sigma \delta \varepsilon \int_0^\infty J_\nu (h\nu)^{\frac{1}{2}} d\nu \qquad (8)$$

proportional to the intensity of the radiation and to ε, and independent of n. This is qualitatively similar to the ordinary radiation pressure, but larger by a factor $\varepsilon(mc^2/h\nu)$, where $\bar{\nu}$ is the average photon frequency.

In the solar neighborhood the average energy density of radiation is 0.01 photons cm^{-3}, Witt and Johnson (1973), and if we assume a mean photon energy of 6 e.v. and a velocity $\tilde{v}$ of about 10^5 cm sec^{-1}, the critical density n_c has a value of $\approx 10^4 \, cm^{-3} \, \varepsilon$. According to the estimate of Gersten et al (1975), $\varepsilon \gtrsim 10^{-3}$ for single atoms. It appears, therefore, that for those diffuse clouds in the solar neighborhood where the grain temperature is low enough, $T_g \lesssim 13$ K, the radiation exerts a force on the grains independent of the radiation intensity and proportional to the gas density. Denser clouds are in the second regime. It is important to remember that the grain jet force at optical depth $\tau = \frac{1}{2} \ln(mc^2/h\bar{\nu}) \sim 10 + \ln \varepsilon$ can be comparable to ordinary radiation pressure at $\tau = 0$.

III. Separation of Gas and Dust.

The grain jet effect is the source of instabilities in the interstellar medium which may lead to the separation of gas and dust and to the separation of grains of differing composition and/or sizes. If we consider a uniform distribution of gas and dust immersed in a homogeneous and isotropic radiation field, and subjected to a perturbation in dust density alone, such a perturbation is unstable. The origin of this instability is physically obvious; a region of increased density attenuates the interstellar radiation field, producing an anisotropy in the intensity. This anisotropy in turn drives the grains toward the region of higher density, amplifying the perturbation.

It is a short calculation to demonstrate that the system described above is unstable to perturbations in the density of the dust component. If we consider the gas and dust to be two fluids interacting only through the drag force on the moving grains, we may note that the grain jet effect does not transfer any net momentum to the gas, except for that given by the radiation pressure on the grains, which we shall neglect. In order to be consistent, we therefore neglect all terms of order 1/c in the equation of motion of the dust. This is equivalent to neglecting the difference in the radiation field as seen in the frame moving with the dust from that seen in the frame fixed to the stars. We may then write the equations of motion for the dust in the form

$$\frac{D\rho}{Dt} + \nabla(\rho\vec{v}) = 0 \tag{9}$$

$$\frac{D\bar{v}}{Dt} = -B\, n_H\, v + A\, H_1/J_0 \tag{10}$$

where the coefficients A and B are given by

$$A = \begin{cases} \epsilon\, \sigma \sqrt{\dfrac{2m}{h\bar{\nu}}}\, J_0; & f = 1 \\ c_s \sigma \sqrt{m\, h\, \bar{\nu}}\; \dfrac{S\, n_H}{m_g}; & f \ll 1 \end{cases} \tag{11}$$

$$B = n_H \sigma c_s \quad . \tag{12}$$

The only photons of interest in our problem are the ultraviolet ones capable of desorbing atoms and/or molecules from the grains; we may thus neglect the emission of the grains and write the transfer equation as

$$\mu \frac{\partial I}{\partial Z} = - kI + k \alpha J + j_* \quad , \tag{13}$$

where α is the albedo of the grains and j_* is the emission by the uniform background of stars. These equations are the starting point for our linear stability analysis.

We introduce a perturbation in density, $\rho_1 = \tilde{\rho} \exp [nt - ikZ]$, along with corresponding perturbations in velocity and radiative intensity. The linearized equations of motion and the equation of transfer become

$$n\rho_1 = i\, k\, v_1\, \rho_0 \tag{14}$$

$$nv_1 = - B\, n_H\, v + A\, H_1/J_0 \tag{15}$$

$$(- i\, k\, \mu + \kappa_0)I_1 + \kappa_0 \frac{\rho_1}{\rho_0} (I_0 - \alpha\, J_0) = \kappa_0\, \alpha\, J_1 \tag{16}$$

Here, $\kappa_0 = k\, \rho_0$ and $\mu = \cos\theta$ as usual. The equation of transfer applies to one frequency, but, since it is linear, we may regard it as applying to all frequencies of interest, with the appropriate opacity.

To simplify the solution of this system for $n(k)$, we introduce the dimensionless variables $\kappa = k/\kappa_0$; $\phi^2 = \frac{1}{\kappa_0 A} n(n - B\, n_H)$; $\xi_1 = H_1/J_0$; $\Pi_1 = J_1/J_0$; $\Psi_1 = I_1/J_0$. If we then eliminate ρ_1 and v_1 by using (14) and (15), we obtain

$$(1 - i\kappa\mu)\Psi_1 + i(1-\alpha)\frac{\kappa}{\phi^2}\,\xi_1 = \alpha\,\Pi_1 \tag{17}$$

Integrating (17) over μ, we have

$$\Psi_1 - i\kappa\;\xi_1 + i(1-\alpha)\frac{\kappa}{\phi^2}\,\xi_1 = \alpha\,\Pi_1\;, \tag{18}$$

while, if we solve (17) for Π_1 and then integrate, we obtain

$$\Pi_1 = \frac{1}{i\gamma}\left[\alpha\,\Pi_1 - i(1-\alpha)\frac{\kappa}{\phi^2}\,\xi_1\right]\;, \tag{19}$$

where

$$\frac{1}{i\gamma} \equiv \frac{1}{i}\,\frac{\tan^{-1}\kappa}{\kappa} = \tfrac{1}{2}\int_{-1}^{1}\frac{d\mu}{1 - i\kappa\mu}\;. \tag{20}$$

Combining the relations (18) and (19), we may eliminate and finally obtain the result

$$\phi^2 = 1 - \alpha + \frac{(1-\alpha)^2}{\alpha - i\gamma}\;. \tag{21}$$

In order to put this result back into dimensional form, we introduce the characteristic time $\tau_0 \equiv N_0^{-1} = (\kappa_0 A)^{-\frac{1}{2}}$, so that

$$n(n + Bn_H) = N_0^2\left\{1 - \alpha + \frac{(1-\alpha)^2}{\alpha - i\gamma}\right\}. \tag{22}$$

For optically thin perturbations ($\kappa >> 1$), we may neglect the second term and simply obtain

$$n(n + B\,n_H) \cong N_0^2\,(1 - \alpha)\ . \tag{23}$$

Notice that, if the albedo of the grains is 1, no photons are absorbed and there is no photodesorption; it is therefore satisfying to note that $\phi^2 = 0$ in this case.

Before discussing the application of this result, we may give a short heuristic derivation of the growth time τ_0. If we consider a perturbation of wavelength λ ($= 2\pi/k$), the change in optical depth in a wavelength is about $\tau_1 = \kappa_0\,\lambda(\delta_\rho/\rho)$. Therefore, the anisotropy of the radiation field will be approximately $\delta = H_1/J_0 \equiv \tau_1$. The grains, which quickly achieve their limiting speed of $A/B\,n_H^{\ \delta}$, are driven through the gas by this anisotropic radiation. In order to increase the magnitude of the perturbation by an amount comparable to δ_ρ, the grains must move a distance $L = \lambda/\delta$. At their limiting velocity, this will take a time $T = L/v = B\,n_H/K_0A$. If we put $B\,n_H >> n$ in (22), corresponding to the fact that the grains achieve their limiting velocity very rapidly, we then have $n^{-1} = (B\,n_H/K_oA)\left\{1 - \alpha + (1 - \alpha)^2/(\alpha - i\gamma)\right\}$. Since the factor in the curly brackets is just a number near one, the growth time derived from this argument is the same as that obtained from the stability analysis.

For the low density regime, which includes most diffuse interstellar clouds, we may estimate N_o by taking $\sigma \sim 10^{-9}$ cm^2; $h\bar{\nu} = 3$ e.v.; $S \simeq 1$; $c_s = 1$ Km/sec. and $m_g = 10^{10}\,m_H$. In this case, $\tau_o = 10^7$ years/n_H. This time is less than the typical lifetime of interstellar clouds; we may therefore expect the grain jet effect to be of importance in the structure of such clouds.

The differences in temperature among grains of different sizes and composition, leading to different thermal evaporation rates (c.f. Eq. (3)), generates a separation among grains with different properties through the grain jet effect. The small grains which radiate inefficiently in the infrared, may become too hot to be driven through the gas by this mechanism, while the larger ones may move rapidly with respect to them and the gas. The quantitative details of this process will depend, of course, on the (unknown) composition of the grains themselves.

In the regions near hot stars, the radiation field itself may be highly anisotropic, and, as Reddish originally suggested (see also Brand, 1975), may drive the dust at high velocities. However, the stellar radiation may heat the grains enough to make the thermal evaporation proceed more rapidly than photodesorption, since the former depends exponentially on the grain temperature. The growth rate n of perturbations in the low density regime of the grain jet effect does not depend on the radiation flux; when the density is high enough, the second limiting case obtains and because of its dependence on the flux of radiation, the instability eventually shuts off. At this stage, however, the dust density should have increased by a very large factor, and our linearized analysis no longer applies.

The second regime of the force is analogous to the radiation pressure. If $\varepsilon = 10^{-4}$, the force due to photodesorption is equal to that due to radiation pressure. For those regions in the second regime of the grain jet effect, using the same values of the parameters as before we find $\tau_o \approx (5 \times 10^2/\varepsilon)e^{\tau}$ years, where τ is the opacity. If $\varepsilon \lesssim 10^{-4}$, the growth time again is less than the typical lifetime of interstellar diffuse clouds.

The second regime also applies to the dynamical effects of the photo-ejection of electrons from the surface (Watson, 1972). The momentum transferred to the grain per photoelectron is smaller than the one transferred by photo-desorbed molecules by a factor $(m_e/m)^{\frac{1}{2}}$, where m is the mass of a typical molecule. This effect is compensated, however, by the very high efficiency expected for the photoelectric effect. Values of this efficiency are very uncertain at present time but according to Watson they range from 0.01 to 0.1. For very small grains, Jura (1975) estimates that the efficiency could be as large as 0.67. For optically thin clouds, therefore, $\tau_o \approx 3 \times 10^5 - 3 \times 10^6$ years.

The point to be emphasized then is that the complete mixture of gas and dust in an isotropic radiation field is an unstable configuration: <u>all diffuse interstellar clouds should have an inhomogeneous dust distribution.</u>

According to the treatment of Spitzer (1968) of the equilibrium electric charge of grains, the neutral subsonic drag force, used in the preceding calculations, dominates over the plasma drag if the ionization degree of the gas is $\lesssim 3 \times 10^{-3}$. For HI clouds we are then justified in neglecting the plasma drag.

In the intercloud medium, however, the reverse is true and in this case we have to multiply (12) by a factor $\simeq 363.\ \chi$, where χ is the degree of ionization. Following the conventional view of the intercloud medium, e.g. Dalgarno and McCray (1972), we adopt $\chi \simeq 0.1$ and $n_e \simeq 0.3$. In this case $\tau_o \simeq 10^9$ years, a value much larger than many other characteristic times of the intercloud medium. Neither the grain jet effect nor the photoelectric effect are able to separate the dust from the gas in the intercloud medium.

IV. Heating of Interstellar Clouds.

At the present time the basic mechanisms responsible for the heating of the interstellar diffuse clouds are not fully understood. In the low density regime of the grain jet effect, which includes diffuse clouds, the rate of ejection of molecules or atoms is independent of the intensity of radiation and the efficiency for photodesorption, being equal simply to the rate of arrival of atoms to the surface. As the adsorption energies are very small compared to the average energy of the photons, every atom that sticks to the surface of the grains is photodesorbed and carries off an energy of several electron volts. This process may, therefore, constitute an important heat source for the interstellar medium, c.f. Silk (1973).

If we assume an average dust to gas ratio of 10^{-12} grains/H atom, the heating rate is approximately

$$H = 5 \cdot 10^{-28} \, n_H^{\,2} (T/80)^{\frac{1}{2}} \text{ ergs s}^{-1}. \tag{24}$$

According to Glassgold and Langer (1974), two of the most important heating mechanisms for clouds are cosmic rays and the photoelectric effect from grains. The ratio of the heating rate due to photodesorption, as given by (24), and the heating rate due to the photoelectric effect, is $(n_H/25)\,(T/80)^{\frac{1}{2}}\,(y/0.1)^{-1}$, where y is the efficiency of the photoelectric effect. It appears then, that for diffuse clouds, photodesorption may be more important than or at least comparable to the photoelectric effect as a source of heating.

Until $n_H < n_c$, the heating is nearly independent of the radiative intensity, and thus is unaffected by shielding due to the grains. The fact that the heating rate has the same density dependence as the cooling rate ought to enhance the thermal stability of the clouds.

Although there is no observational evidence for H_2 molecules with kinetic energies as high as 6 e.v. in the diffuse clouds observed by the Copernicus satellite (Spitzer and Jenkins, 1975), the thermalization time for such molecules is sufficiently short if $n > 10\ cm^{-3}$ that this is compatible with grain jet heating for such clouds. If the velocities of ejection are smaller it is doubtful that this process is an important source of heating; its dynamical effects, however, can still be substantial and this is the point we want to emphasize in this note. Equating the heating rate given by (24) to the cooling rate due to CII, the main cooling mechanism in diffuse clouds (Dalgarno and McCray, 1972), we estimate a cloud temperature of $\sim 90^{\circ}$ K, independent of the density.

V. Discussion.

On the basis of the preceding results it is tempting to speculate about typical consequences. The clumps of dust formed by the instability may be sites of very rapid formation of molecular hydrogen if they become optically thick. Although it is not possible to calculate the local enhancement of the dust to gas ratio from the linear theory of instability, simple numerical estimates indicate that increases on this ratio of very large factors are to be expected. Diffuse clouds should thus have an inhomogeneous molecular distribution.

Because of the heating mechanism, the gas temperature inside these clumps of high dust density would increase; observations of molecular hydrogen, such as those by the Copernicus satellite, would tend then to sample the regions of highest temperature in a cloud. The temperature inhomogeneities would also generate motions of the gas to equalize the pressure differences.

In summary, we may conclude that the grain jet effect may have important consequences in all the neutral parts of the interstellar medium except the diffuse intercloud regions. In the moderate density regions ($n < n_c$), the force per grain is proportional to the gas density and the radiation anisotropy, and independent of the radiative intensity. The contribution of photodesorption to the energy balance of the interstellar gas may be important in this regime. In regions of high density, the force on the grains is independent of the density and proportional to the radiative flux. In such regions, it exceeds the radiation pressure by large factors.

If the parameters discussed above have values which lie well within the range of those usually assumed, this process may be an important source of heat in dense HI regions, the origin of new instabilities whose consequence for the formation of dense clouds, molecules and stars must be explored, and a mechanism for producing large variations in the ratio of gas to dust in the galaxy.

As with all mechanisms which depend on properties of the interstellar grains, many uncertainties still remain in the basic physics. We hope that the discussions in this paper will be a stimulus for further laboratory studies on photodesorption from solid surfaces as well as on research on its astrophysical consequences.

We are greatly indebted to a referee for his constructive comments on the original version and to Profs. A. E. Glassgold and F. Kahn for interesting discussions. This research has been supported in part by grant NGR-33-016-196 from the National Aeronautics and Space Administration.

REFERENCES

Brand, P.W.J.L., 1975, Astr. and Space Sci. 33, 231.

Dalgarno, A. and McCray, R., 1972, Ann. Rev. Astr. and Ap. 10, 375.

Fabel, G. W., Cox, S. M. and Lichtman, D., 1973, Surface Sci. 40, 571.

Gersten, J. I., Janow, R. and Tzoar, N., 1975, Phys. Rev. B 11, 1267.

Glassgold, A. E. and Langer, W. D., 1974, Ap. J. 193, 73.

Greenberg, L. T., 1973, Interstellar Dust and Related Topics, Greenberg and Van de Hulst, Eds. (D. Reidel), 413.

Hollenbach, D. and Salpeter, E. E., 1971, Ap. J. 163, 155.

Jura, M., 1975, preprint.

Leung, C. M., 1975, Ap. J. 199, 340.

Reddish, V., 1971, Nature 232, 40.

Silk, J., 1973, PASP 85, 704.

Spitzer, L., 1968, Diffuse Matter in Space (Interscience).

Spitzer, L. and Jenkins, E. B., 1975, Ann. Rev. Astr. and Ap., in press.

Watson, W. D., 1972, Ap. J. 176, 103.

Watson, W. D. and Salpeter, E. E., 1972, Ap. J. 174, 321.

Witt, A. N. and Johnson, M. W., 1973, Ap. J. 181, 363.