Nova Dust Nucleation: Kinetics and Photodissociation

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Abstract

Dust is observed to form in nova ejecta. The grain temperature is determined by the diluted nova radiation field rather than the gas kinetic temperature, making classical nucleation theory inapplicable. We used kinetic equations to calculate the growth of carbon nuclei in these ejecta. For expected values of the parameters too many clusters grew, despite the small sticking probability of atoms to small clusters, and the clusters only reached radii of about 100Å when the carbon vapor was depleted. We then included the effects of cluster photodissociation by ultraviolet radiation from the nova. This suppresses nucleation, but too well, and no grains form at all. Finally we suggest that a few growing carbon nuclei may be protected from photodissociation by a sacrificial surface layer of hydrogen.

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I. Introduction

Novae can be divided into two general classes: fast novae which fade more rapidly than 0.1 magnitude day$^{-1}$ and have higher average luminosities, and moderate speed and slow novae which fade less rapidly than 0.1 magnitude day$^{-1}$ and have lower average luminosities. The decline in the visual light curve for moderate speed novae is accompanied by an increase in infrared brightness, with the infrared luminosity matching the initial bolometric luminosity of the nova. These observations have been interpreted as resulting from the condensation of dust grains in an expanding shell of ejecta. This generally accepted model is built on the assumption that visible and ultraviolet radiation from the nova is absorbed by the dust, and is reradiated in the infrared. Fast novae, in contrast, do not show this evidence for dust production.

Because infrared spectral signatures, such as the characteristic 10\(\mu\) silicate line, have not been generally observed from these shells, it is believed that the dust is largely carbonaceous. Further analysis (Ney and Hatfield 1978; Gehrz, et al. 1980ab; Bode and Evans 1982; Mitchell, et al. 1983) led to typical dust grain diameters of a few tenths of a micron or, in some models, as large as a few microns, and to visible optical depths $\sim$ 3. However, extensive modeling of grain nucleation and growth, for example by Yamamoto and Hasegawa (1977) and Draine and Salpeter (1977), has not been able to reproduce the grain size distributions that the emission models appear to require.

Donn and Nuth (1985) have given several reasons for why classical nucleation theory is inapplicable in many astrophysical environments. Classical nucleation theory uses macroscopic parameters like surface tension to calculate the free energy of tiny clusters, and assumes the gas temperature and the cluster vibrational temperature are equal. In fact, cluster temperatures are usually controlled by the radiation field, and do not equal the gas temperature. The theory assumes in its derivation that very small clusters are present in equilibrium concentrations. This is incorrect around novae, because the densities are so small and conditions change so quickly. Carbon monomers may have a number density of
$10^7$ cm$^{-3}$ and a thermal velocity of $4 \times 10^5$ cm sec$^{-1}$, a collision cross-section of $4 \times 10^{-16}$ cm$^2$, and a mean time between collisions of $\approx 600$ seconds. With a sticking coefficient of $\approx 10^{-6}$ (Draine 1979), the average carbon atom would become part of a carbon dimer in $6 \times 10^8$ seconds. The dust formation process takes place within $3 \times 10^6$ seconds, so dimers will not be in equilibrium with monomers. Classical nucleation theory is therefore inapplicable to nova ejecta.

Because nucleation theory is inapplicable and fails to account for the observations, we have performed kinetic calculations of the growth of small condensation nuclei. Calculations using theoretical sticking probabilities on small clusters led to grain sizes which were far too small to explain the infrared and visible data. Too many clusters grow, and the carbon vapor is depleted before they reach the empirically inferred size. We then realized that the presence of a significant ultraviolet flux would efficiently photodissociate small grains. The absorption of even a single photon raises the temperature of a small cluster sufficiently to induce evaporation of a carbon atom, a process analogous to photodissociation of a molecule. This offered hope of avoiding the excess of condensation nuclei. Our calculations showed that this destruction mechanism is, indeed, too efficient, and that it effectively prevents the formation of any grains at all. It is possible that processes we have not included, such as the effects of hydrogenation of the grain surfaces, are crucial, or that grains form heterogeneously on pre-existing nuclei.

II. The Nucleation Model

In our study of nucleation around novae we integrated the following set of 60 equations using the IMSL routine DGEAR: A volume $V$ contains $N_g$ clusters of size $g$, each cluster containing $f(g)$ carbon atoms. We chose $f(g) = g$ for $1 \leq g \leq 20$, $f(g) = 20 \times 2^{(g-20)/5}$ for $21 \leq g \leq 24$, and $f(g) = 20 \times 2^{g-24}$ for $g \geq 25$. We describe clusters of carbon atoms from monomers up to grains with $20 \times 2^{36} \approx 1.37 \times 10^{12}$ atoms, or grains with a radius
\( \approx 1\mu \). The evolution equations for \( 3 \leq g \leq 60 \) are

\[
d\frac{N_g}{dt} = \frac{v_{1,g-1}A_{g-1}N_{g-1}N_1}{(f(g) - f(g-1))V(t)} - \frac{v_{1,g}A_gN_gN_1}{(f(g+1) - f(g))V(t)} + \frac{v_{2,g-1}A_{g-1}N_{g-1}N_2}{(f(g) - f(g-1))V(t)} - \frac{v_{2,g}A_gN_gN_2}{(f(g) - f(g-1))V(t)} - \frac{\beta_gN_g}{f(g) - f(g-1)} + \frac{\beta_{g+1}N_{g+1}}{f(g+1) - f(g)}. \tag{1}
\]

For \( g = 2 \)

\[
d\frac{N_2}{dt} = -\sum_{g=3}^{60} \frac{v_{2,g}A_gN_gN_2}{(f(g) - f(g-1))V(t)} - \frac{2v_{2,2}A_2N_2^2}{V(t)} - \frac{\alpha_2v_{1,2}A_2N_2N_1}{V(t)} + \beta_3N_3 - \beta_2N_2 + \frac{\alpha_1v_{1,1}A_1N_1^2}{V(t)}. \tag{2}
\]

For \( g = 1 \)

\[
d\frac{N_1}{dt} = -\sum_{g=2}^{60} f(g) \frac{dN_g}{dt}. \tag{3}
\]

Here \( A_g \) and \( \alpha_g \) are the cross-section and sticking coefficient respectively for an atom colliding with a \( g \)-cluster. For collisions of dimers with larger clusters the sticking coefficient is assumed to be 1; the dimer is assumed to split and the excess energy to be carried off by one atom which does not stick to the cluster, while the cross-section is taken to be the same as that for atoms. The evaporation coefficient for \( g \)-clusters is \( \beta_g \) and includes both ordinary thermal evaporation and the evaporation which takes place immediately after a photon is absorbed. \( V(t) \) is the (arbitrary) volume of the volume element we examine and is initially equal to 1 cm\(^3\) when we start the calculation. \( v_{1,g} \) is the mean relative velocity of a \( g \)-cluster and a monomer, while \( v_{2,g} \) is the mean relative velocity of a \( g \)-cluster and a dimer.

The collision cross-section is taken as \( f(g)^{2/3}\pi a_0^2 \), where \( a_0 \) is the radius of a carbon atom. The sticking coefficient was calculated as suggested by Williams (1972), Donn et al. (1981), and Freed et al. (1982): When two atoms or an atom and a cluster collide they form a new, larger cluster with more than enough energy to dissociate. The excess energy must be disposed of somehow or the new transient cluster will break up. In terrestrial
environments the density of atoms or molecules is high enough that a third body may collide with the excited cluster and carry off the excess energy. Around novae such three body reactions are very unlikely. Collisions with hydrogen may occur on a time scale of \((nvA)^{-1}\) where \(n \approx 10^9\), \(v \approx 10^6\) cm sec\(^{-1}\), and \(A \approx 10^{-15}\) cm\(^2\); the time scale is \(\sim 1\) second. An excited dimer will only last about \(10^{-13}\) second, so that only one collision in \(10^{13}\) will be stabilized by a third body.

In such low density environments stabilization is more likely to occur by the emission of a photon which can carry off the excess energy. Such reactions are called radiative association reactions. The sticking coefficient, called the probability of radiative stabilization by Donn, et al. (1981), is

\[
\alpha = \frac{\tau_r^{-1}}{\tau_r^{-1} + \tau_d^{-1}},
\]

where \(\tau_r\) is the radiative lifetime of the excited cluster and \(\tau_d\) is its dissociative lifetime. We need to determine these two quantities.

The radiative lifetime \(\tau_r\) is \(\sim 10^{-7}\) second if emission takes place by an allowed electronic transition, and is \(\sim 10^{-2}\) second if cooling takes place by an allowed vibrational transition. Clearly sticking is more likely if electronic transitions occur. For the formation of dimers \(\tau_d\) is equal to a vibrational period, or \(\sim 10^{-13}\) second. Then for \(C + C \rightarrow C_2\), the sticking coefficient is either \(\sim 10^{-6}\) or \(\sim 10^{-11}\), depending on which sort of emission process takes place. The empirical sticking coefficient of \(\sim 10^{-6}\) (Draine 1979) implies that an electronic transition occurs (also, vibrational transitions are forbidden for \(C_2\) and their actual rate is much slower than \(10^2\) sec\(^{-1}\)). We generalize from dimers by using the radiative lifetime of excited electronic states in calculating all sticking coefficients. If this is incorrect it will overestimate the sticking coefficient.

The other time scale is \(\tau_d\), the average lifetime of the excited cluster before it breaks up. We estimate this using RRK theory (Johnston 1966, Duley and Williams 1984), which assumes that the cluster is a collection of harmonic oscillators of some frequency \(\nu\), which for carbon clusters we take to be \(4.5 \times 10^{13}\) sec\(^{-1}\) (Freed et al. 1982). One particular bond
(or vibrational normal mode) is assumed to be more likely than the others to break up, and if this mode accumulates enough energy dissociation occurs. The internal energy of the cluster is distributed among the various oscillators in packets of magnitude $h\nu$. The dissociation energy $E_d = Mh\nu$. The total energy of the cluster is $E = Jh\nu$. The number of oscillators $S$ is equal to the number of vibrational modes of the cluster: $S = 3N - 6$ for $N \geq 3$ (we take these molecules to be nonlinear, although in fact $C_3$ is linear) and $S = 1$ for dimers. The number of ways to distribute $J$ identical packets or quanta of vibrational energy among $S$ modes is $\frac{(J+S-1)!}{J!(S-1)!}$. For dissociation to occur at least $M$ quanta must be in one particular vibrational mode. Then the number of ways to distribute the remaining $J - M$ quanta over the $S$ oscillators is $\frac{(J-M+S-1)!}{(J-M)!(S-1)!}$. The fraction of of the total number of distributions which will lead to dissociation is $\frac{(J-M+S-1)!}{(J+S-1)!(J-M)!}$. If rearrangements of quanta occur at a rate $\nu$, the dissociation rate is

$$\tau_d^{-1} = \nu \frac{(J-M+S-1)!J!}{(J+S-1)!(J-M)!}. \quad (5)$$

If there are $q$ modes which lead to dissociation, we multiply this expression by $q$. For dimers, $S = 1$ and $\tau_d^{-1} = \nu$.

To calculate the sticking coefficient we write the mean energy of a cluster which has just been hit by a gas-phase monomer with the mean kinetic energy $1.5k_BT_{gas}$:

$$E = Sk_BT_{grain} + 1.5k_BT_{gas} + E_d, \quad (6)$$

where $T_{grain}$ is the grain temperature, $T_{gas}$ is the gas temperature, $k_B$ is Boltzmann’s constant, and $E_d$ is the dissociation energy which we take to be 6 eV. We then have $J = E/h\nu$ and $M = E_d/h\nu$, and can calculate $\tau_d^{-1}$. The sticking coefficient $\alpha$ is shown in Table 1.

The sticking coefficient falls short of 1, even for very large clusters, because then $J \approx S(k_BT_{grain}/h\nu) \gg M$ and

$$\tau_d^{-1} \to \nu \left( \frac{J}{J+S-1} \right)^M \to \nu \left( \frac{k_BT_{grain}}{k_BT_{grain} + h\nu} \right)^M. \quad (7)$$
This is much less than $\nu$, but it is not zero.

The equation for relative velocity is taken from Draine and Salpeter (1979), who give an expression for the drag force $F_{\text{drag}}(v)$ on a neutral grain moving through a gas. We balance this velocity dependent drag force against the force of radiation pressure:

$$F_{\text{rad}} = \frac{L}{4\pi r_C^2} Q \pi a^2,$$

where $a$ is the grain radius, $L$ the nova’s luminosity, and we take the effective momentum transfer coefficient $Q = \min(10^5 a, 1)$. Equating $F_{\text{drag}}(v_{gr}) = F_{\text{rad}}$ gives $v_{gr}$, the velocity of a grain through the gas. We can then calculate the flux of carbon atoms hitting a grain using the equations derived by Aannestad and given by Shull (1978). The result is

$$F = \left(\frac{\pi}{2}\right)^{1/2} a^2 n_C v_{th} \frac{v_{gr}}{v_{gr}} \left(2\pi v_{th}^2 \left(1 + \frac{v_{gr}^2}{v_{th}^2}\right) \text{erf} \left(\frac{v_{gr}}{2^{1/2} v_{th}}\right) + 2v_{gr} \exp \left(-\frac{v_{gr}^2}{2v_{th}^2}\right)\right),$$

where $v_{th} \equiv (k_B T_{\text{gas}}/m_C)^{1/2}$ is the atomic thermal velocity and $n_C$ is the density of carbon atoms. When $v_{gr} \ll v_{th}$ this reduces to

$$F = \pi a^2 n_C v_{gr},$$

as it must.

For our system of grain-growth equations (1)–(3), we define the relative velocity $v$ so that $n_C v \pi a^2$ gives the flux. Then

$$v = \frac{F}{\pi a^2 n_C}.$$

We calculate this velocity for collisions between monomers, dimers and larger clusters. For the relative velocities of two monomers or monomers and dimers, we used $v = (8/\pi)^{1/2} v_{th}$, which is not strictly correct, but adequate given the uncertainty in the reaction cross-section. For dimer-dimer collisions we similarly used $v = (4/\pi)^{1/2} v_{th}$.

We split the evaporation rate into two terms, one giving the evaporation rate at the cluster temperature $T_m$ corresponding to the mean heat flux, and the other giving
the evaporation rate of the cluster as a result of absorption of single ultraviolet photons (photodissociation). This is a somewhat artificial distinction, because the temperature of the cluster is determined by a balance of photon absorption and infrared emission. For very small clusters the temperature fluctuations are much larger than the average temperature, while for clusters composed of many atoms the absorption of a single ultraviolet photon has little effect on the temperature. The method outlined in the Appendix treats dissociation in a unified way. Here we describe a simpler approach.

$T_m$ is determined by the equation

$$T_m = \left( \frac{L}{16\pi\sigma r^2 Q_{UV}} \right)^{1/4}, \quad (12)$$

where the ratio of emissivities $Q_{UV}/Q_{IR} = (T_*/T_m)^{1.65}$ (Clayton and Wickramasinghe 1976), $\sigma$ is the Stefan-Boltzmann constant, and $T_*$ is the color temperature of the nova. We adopt $T_* = 8000^\circ$K. Then the evaporation rate of a spherical cluster is

$$\beta = 4\pi a^2 \alpha \left( \frac{P_s}{(2\pi m_C k_B T_m)^{1/2}} \right) \exp \left( \frac{2\gamma \Omega}{a k_B T_m} \right), \quad (13)$$

where the exponential term comes from the Gibbs correction to the vapor pressure, $P_s = 1.66 \times 10^{14} \exp(-88880^\circ$K/$T_m)$ (Lefevre 1979) is the saturation vapor pressure over a flat surface, $\alpha$ is the sticking coefficient, $\Omega$ is the volume of an atom in the condensed phase, and $\gamma = 1000$ erg cm$^{-2}$ is the surface energy per unit area. For large grains the mean evaporation rate (13) exceeds that resulting from fluctuations.

For small clusters fluctuations dominate. A cluster with 4 atoms has 6 degrees of freedom. A 6 eV photon will give an average of 1 eV of energy to each degree of freedom, increasing its temperature by $11000^\circ$K. We then must calculate the probability of dissociation occurring while this high temperature lasts—the photodissociation rate.

III. Photodissociation

In order to calculate the photodissociation rate we must calculate the probability that a cluster reradiates its absorbed energy before it dissociates. If the average radiating time
is $\tau_r$ and the average dissociating time is $\tau_d$, then the chance of its dissociating before it loses energy by radiation is $\tau_d^{-1}/(\tau_d^{-1} + \tau_r^{-1})$. The dissociation rate from RRK theory is given by (5). This rate has its smallest value if $E = E_d (J = M)$. Then

$$\tau_d^{-1} = \nu \frac{(S - 1)!M!}{(M + S - 1)!}. \quad (14)$$

The larger the cluster the more degrees of freedom it has, and the less likely it is to dissociate. The probability of breakup is

$$P = \frac{\tau_d^{-1}}{\tau_r^{-1} + \tau_d^{-1}}. \quad (15)$$

Omont (1986) describes what happens when a polycyclic aromatic hydrocarbon molecule, which resembles a small carbon cluster, absorbs an ultraviolet photon. If the energy is great enough for photoionization, ionization can occur with a probability $> 0.5$. Photodissociation can also occur immediately, if the excited electronic state is vibrationally unbound. If neither of these events occurs, the energy generally is thermalized as vibrational energy (reradiation by an electronic transition is possible, and is accounted for by use of an empirical absorption cross-section). The cluster then cools by vibrational infrared radiation or by evaporation.

We show in Table 2 dissociation times $\tau_d$ of pure carbon clusters, assuming that their internal energy exactly equals the dissociation energy, taking $\nu = 4.5 \times 10^{13} \text{ sec}^{-1}$, $M = 30$, and $S = 3N - 6$, and using (14). Clusters of $N < 8$ will readily photodissociate. Unless the ultraviolet spectrum is very hard, only one carbon atom will be lost for each absorbed photon with $h\nu_{UV} > E_d$.

It is important to know at what wavelengths the clusters in question can absorb. A well-known paper by Platt (1956) estimated the longest wavelength which could be absorbed by a molecule with unsaturated energy bands. Platt treats the molecule as a rectangular box with dimensions equal to the dimensions of the molecule. Each atom is assumed to contribute one electron and the lowest energy levels are filled up. The amount
of energy it takes to raise an electron from the highest filled energy level to the next one above it gives, in this model, the lowest energy photon which the molecule can absorb. For linear molecules the wavelength of the photon is about 400 times the length of the molecule. Platt found that experimental data for many different organic molecules agreed with this prediction to within a factor of two. Thus, molecules about 10 Å long are just barely able to absorb in the visible region. Certainly even very small molecules will be able to absorb ultraviolet light. We will ignore absorption by visible photons, which have energies between 2 and 4 eV. Two or more successive absorptions of these photons would provide enough energy to break loose a carbon atom, but unless the absorption rate is very high the molecule will reradiate much of its energy in the interval between absorptions; hence we ignore this route to photodissociation.

Having decided only to include ultraviolet photons in our simulation, we must assign the small clusters an absorption cross-section. From Lee (1984), it appears that for many small molecules the absorption cross-section is roughly $10^{-18}$ cm$^2$. Omont (1986) gives a value of $\approx 5 \times 10^{-19}$ cm$^2$ per carbon atom for the absorption cross-section of polycyclic aromatic hydrocarbons at visible wavelengths; for the ultraviolet wavelengths the cross-sections are an order of magnitude larger. We adopted a value of $1.5 \times 10^{-18}$ cm$^2$ atom$^{-1}$.

IV. Results

We first describe calculations without photodissociation. The evaporation rate was given by (13), using the mean temperature (12). In one calculation we assumed that $10^{27}$ gm of carbon atoms were present and spread evenly throughout a shell of volume $4\pi r^3/10$; i.e., at a given instant there was a shell of radius $r = vt$ and of thickness $\approx r/10$. We began the calculation at a time after the nova outburst $t = 10^6$ sec, and the expulsion velocity $v = 750$ km sec$^{-1}$. Nucleation begins around $t = 2.5 \times 10^6$ sec. By $t = 3.4 \times 10^6$ sec less than 1% of the carbon remains monomeric. The distribution of grain sizes is shown by the solid line in Figure 1. Most of the mass is in grains with $a \approx 100$ Å,
in disagreement with the infrared data. The visible optical depth (calculated using the opacity of Clayton and Wickramasinghe 1976) is 90, much greater than observed. If the shell were 30 times thinner (but of the same density) and the carbon mass were 30 times less, the optical depth would be 3, as observed for Nova Vul 1976, Nova Ser 1978, and Nov Ser 1970. Such a thin shell, equal to $r/300 \sim 10^{12}$ cm in thickness, is unlikely to form and would not remain thin for a period greater than $\Delta r/2v_s$ (Ennis, et al. 1977), where $\Delta r$ is its thickness and $v_s$ is its sound speed.

In another calculation the carbon mass was arbitrarily increased to $10^{28}$ gm in order to see if larger grains would result. Most of the carbon formed dust by $t = 2.6 \times 10^6$ sec. The distribution of grain sizes is shown by the dashed line in Figure 1. Most of the carbon is in grains with radii $a \approx 0.15 \mu$, roughly 10 times greater than before. This is perhaps large enough to agree with the infrared data. However, the visible optical depth was $\approx 700$, which is grossly excessive. To keep the same density of carbon, and therefore the same grain size, but to have the right visible optical depth, we must reduce the mass and shell thickness by a factor of roughly 300, so that $\Delta r = r/3000 = 10^{11}$ cm. This is far too small. We have the same problem that classical nucleation theory has—to produce large grains ($a > 0.1 \mu$) we require a very high density of carbon. The resulting visible optical depth is too great.

It might be that high energy photons keep the number of nucleation sites well under the number predicted by our model. If so, the carbon might accrete onto relatively few nucleation sites to form larger grains. This is possible, if the nucleation sites are sufficiently rare, because once a nucleus overcomes the barrier at $N \leq 7$ it grows rapidly. It was with hopes that this might prove to be the case that we included photodissociation in the calculation. We assumed a blackbody nova spectrum with $L = 1.25 \times 10^{38}$ erg sec$^{-1}$ and $T_s = 8000^\circ$K. All photons with $6.0$ eV $< h\nu_{UV} < 13.6$ eV were assumed to contribute to the photodissociation rate, although in using RRK theory (5) to describe photodissociation we took each to have an energy $h\nu_{UV} = 6.5$ eV, only slightly above the
assumed photodissociation threshold $E_d = 6$ eV.

In separate calculations we took the radiative decay time of the excited state as $10^{-7}$ second, characteristic of an excited electronic state, and $10^{-2}$ second, characteristic of excited vibrational states. The longer radiative time is appropriate (Omont 1986) because the RRK theory applies to vibrationally excited molecules. The choice of radiative times for the molecule shifts the dividing line between small clusters which are almost certain to dissociate, and large clusters which are are almost certain not to dissociate.

We assumed that $10^{28}$ gm of carbon (corresponding to $10^{-4} M_\odot$ of ejecta, of which 5% by mass is carbon) were spread throughout a shell of volume $4\pi r^3/10$, which is expanding at a speed of 750 km/sec. The calculation was begun at the time $t = 10^6$ sec, and resembled our previous calculation except for the inclusion of photodissociation. The results are dramatically different in that no dust forms for either value of the radiative decay time. Virtually no dimers form; with $\tau_r = 10^{-2}$ second, at the calculation’s end the ratio of dimers to monomers is $6 \times 10^{-10}$. Essentially all the carbon remains as monatomic vapor.

The bottleneck to grain formation occurs at very small molecules: $C_2$, $C_3$, $C_4$, etc., for which absorption of a 6.5 eV photon virtually guarantees dissociation, since $\tau_d$ is less than $10^{-7}$ second for these tiny clusters. It does not matter whether $\tau_r$ is $10^{-7}$ second or $10^{-2}$ second, since the clusters for which $\tau_d$ is comparable even to $10^{-7}$ second form in only extremely tiny numbers, although the larger (and physically expected) $\tau_r$ worsens the problem.

We tried again with $\tau_r = 10^{-2}$ second for the clusters and $10^{28}$ gm of carbon spread throughout a volume of $4\pi r^3/100$. The radiation rate is realistic, but such a thin shell, with thickness $\approx r/100$, is not. It does give a very high carbon density. Nonetheless, no nucleation occurs. The ratio of dimers to monomers is increased tenfold, as one expects. Still the number of dimers is only $6 \times 10^{-9}$ of the number of monomers and the number of trimers is $10^7$ times rarer still.

We attempted one more run, with a carbon mass of $10^{28}$ gm, a shell volume of $4\pi r^3/10$,
an expulsion velocity of 750 km sec$^{-1}$, and $\tau_r = 10^{-2}$ second. In this last case, for our photodissociation rate we used the same equation (13) that we used for the evaporation rate, with $T$, the “temperature” of the molecule set equal to $E_p/S$. Here $E_p$ is the energy of the photon, set equal to 6.5 eV, and $S$ is the number of degrees of freedom, and equals $3N-6$ (or 1, for dimers). The formula has the advantage of being consistent with the equation we use to calculate the evaporation rate of the clusters at a constant temperature. It does not really matter, for the photodissociation rate for the small clusters is identical to the rate one finds with RRK theory. It merely equals the photon absorption rate. The results are then of course the same: no nucleation and no dust.

No dust forms around novae in our nucleation calculations because ultraviolet photons disrupt small clusters of carbon atoms before they can grow. The exact method we use to determine the rate of dissociation of an excited cluster is not important, because the photodissociation rate of small clusters equals their photon absorption rate.

V. Discussion

A simple model can be constructed to show how the photodissociation rate affects nucleation. Let the equation for the number density of $i$-mers $n_i$ be written as

$$\frac{dn_i}{dt} = R_{i-1}n_{i-1}n_1 - k_in_i - R_in_in_1 + k_{i+1}n_{i+1},$$  \hspace{1cm} (16)$$

where $R_i = v\sigma_i\alpha_i$. For simplicity we set $v\sigma_i = 10^{-10}$ cm$^3$ sec$^{-1}$ for all small clusters. For the sticking coefficient we take $\alpha_i = \min(10^{-7+i}, 1)$. This approximation is slightly more optimistic than that used by Donn, et al. (1981) and than the numbers given by the RRK theory. For $k_i$ we assume a photodissociation cross-section of $10^{-18}$–$10^{-17}$ cm$^2$. The number of photons with $E_p > 6$ eV emitted by a nova with $L = 1.25 \times 10^{38}$ erg sec$^{-1}$ depends on its effective temperature and is in the range $10^{47}$–$10^{49}$ sec$^{-1}$. At a distance of $3 \times 10^{14}$ cm the flux will be $10^{17}$–$10^{19}$ photons cm$^{-2}$. Then the $k_i$ are in the range $10^{-1}$–$10^2$ sec$^{-1}$. 

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For $n_1$ we have a range to pick from. Let the carbon mass be $10^{27} - 10^{28}$ gm. Then, given a shell of volume $4\pi r^3/10$, the carbon density at $r = 3 \times 10^{14}$ cm is $1.6 \times 10^6 - 1.6 \times 10^7$ cm$^{-3}$. A $10^{28}$ gm shell of thickness $r/100$ with clusters forming at $r = 10^{14}$ cm gives $n_1 = 5 \times 10^9$ cm$^{-3}$. We assume (optimistically) that $C_6$ is immune to photodissociation (hydrogenated $C_6$ may be the stable benzene molecule); once a cluster reaches $C_6$ it forms a nucleation site. In reality it must overcome more hurdles and grow still larger to be completely safe from destruction from photodissociation. Then the nucleation rate must be less than $R_5 n_5 n_1$, the formation rate of $C_6$.

To obtain useful analytic results from these equations we must assume that a quasi-equilibrium state is reached so that $\frac{dn_i}{dt} = 0$. This is not a reasonable assumption to make when there is no photodissociation, as Yamamoto and Nishida (1977) demonstrated. It is reasonable in our case: In equilibrium $n_{i+1} \ll n_i \ll n_{i-1}$. Then, neglecting the last two terms on the right hand side of (16),

$$n_i \approx \frac{R_{i-1} n_{i-1} n_1}{k_i}. \tag{17}$$

Later, when we substitute numbers for $n_1$, $R_{i-1}$, and $k_i$, we will see that our assumption that $n_i/n_{i-1} \ll 1$ is justified. The equilibrium value is approached on a time scale $O(k_i^{-1})$. Now $k_i$ is between 0.1 and 100 sec$^{-1}$, so that it takes no more than $\sim 10$ seconds to reach equilibrium. The expansion time scale (the time in which a significant change in the volume of the ejecta takes place) is much greater, justifying the use of the quasiequilibrium distribution.

We can rewrite the equilibrium distribution as

$$n_i \approx \frac{R_{i-1} R_{i-2} \ldots R_1 n_1^i}{k_1 k_2 k_3 \ldots k_i}. \tag{18}$$

Then the nucleation rate $J < R_5 n_5 n_1$, or

$$J < \frac{R_5 R_4 R_3 R_2 R_1 n_1^6}{k_5 k_4 k_3 k_2}. \tag{19}$$
Our previous estimate for the $\alpha_i$ gives $R_i = 10^{-17+i} \text{ cm}^3 \text{ sec}^{-1}$ ($i \leq 7$), and we set $k_i = k$, so that all of the small clusters have the same photodissociation rate (roughly justified because the photodissociation probability is nearly unity). Then

$$n_i \approx 10^{-17(i-1)} 10^{(i^2-i)/2} n_1^i \text{ cm}^3 \text{ sec}^{-(i-1)},$$

or

$$n_5 \approx 10^{-58} \text{ cm}^{12} \text{ sec}^{-4} n_1^5 / k^4$$

and

$$J < 10^{-70} \text{ cm}^{15} \text{ sec}^{-5} n_1^6 / k^4.$$  

In $\Delta t = 10^6$ sec, there will be $J\Delta t$ nucleation sites produced per unit volume, or about $\xi \equiv J\Delta t/n_1$ per carbon atom. This is the maximum possible number of nucleation sites that could be produced during the nova eruption under these assumptions. Then $\xi < 10^{-64} \text{ cm}^{15} \text{ sec}^{-4} n_1^5 / k^4$.

If at least 10% of the carbon condenses, we require $\xi \sim 10^{-10}$ to obtain 0.1$\mu$ grains, or $\xi \sim 10^{-13}$ to form 1$\mu$ grains. For plausible values such as $n_1 = 10^8 \text{ cm}^{-3}$ and $k = 1 \text{ sec}^{-1}$ we find $\xi < 10^{-24}$ and no dust forms, just as we found from numerical integration of the equations (1)–(3). If $n = 10^9 \text{ cm}^{-3}$ and $k = 0.1 \text{ sec}^{-1}$, which are very optimistic, $\xi < 10^{-15}$ and very little dust will form. Only for extreme values $n_1 = 10^{10} \text{ cm}^{-3}$ and $k = 0.1 \text{ sec}^{-1}$ can $\xi$ be as large as $10^{-10}$. This corresponds to $10^{27}–10^{28}$ gm of carbon compressed into a volume of $5 \times 10^{39}–5 \times 10^{40}$ cm$^3$, i.e., a shell of radius $10^{14}$ cm and thickness $10^{11}–10^{12}$ cm. Alternatively, one could assume a clump of gas $\approx 10^{13}$ cm in radius which forms the dust. This would require very asymmetric mass ejection. Mitchell et al. (1985) rejected this possibility. Even if we have extremely high densities, the photodissociation rate must be $< 10^{-1} \text{ sec}^{-1}$ for dust to form. For a black-body source of $L = 1.25 \times 10^{38}$ erg sec$^{-1}$ at $r = 3 \times 10^{14}$ cm and a photodissociation cross-section of $10^{-18}$ cm$^2$ this would require $T_\ast < 7000^\circ\text{K}$.

VI. Conclusions
In our nucleation calculation we showed that it is difficult to understand how nucleation can take place at all. The first obstacle is photoionization; if the radiating temperature of the nova is greater than 12000 °K it is likely that the carbon is entirely ionized and no nucleation can take place. Assuming that in the early stages of the eruption the nova is cooler than this, there are still a great many photons in the 6-11 eV range, below the carbon ionization threshold but still capable of photodissociating small molecules and clusters. They prevent nucleation.

There may be ways around the photodissociation problem. Some of the effects of hydrocarbon chemistry on dust nucleation around novae were discussed by Rawlings and Williams (1989). Omont (1986) pointed out that excited polycyclic aromatic hydrocarbons are more likely to lose a hydrogen atom than a carbon atom, and this will also be true for grain nuclei with any hydrogen attached. Hydrogen is more loosely bound and it is thus more likely to accumulate the necessary amount of energy to break away. The remaining energy will usually be too low for any remaining atoms to escape. Using RRK theory (5) we see that a hydrogen atom is \( \zeta \) times more likely to break off, where

\[
\zeta = \frac{(J - H + S - 1)! (J - C)!}{(J - H)!(J - C + S - 1)!},
\]

\( Hh\nu \) is the binding energy of a hydrogen atom, and \( Ch\nu \) is the binding energy of a carbon atom. If, for instance, \( C = 30, H = 20, J = 30, \) and \( S = 6, \) a given hydrogen atom is 3000 times more likely to be lost than a given carbon atom. This ratio drops steeply with increasing \( J \) (for \( J = 35 \) it is only 60) and decreasing \( S, \) but a large effect is possible.

We neglected hydrogen in our calculations of grain growth, but its number density in the nova debris is typically 100–1000 times that of carbon, and 300–3000 hydrogen atoms collide with a cluster for every carbon atom which does so. Any carbon cluster that does form will likely be hydrogenated on the outside. Collision with a carbon atom will exothermically replace a hydrogen by the carbon; the presence of hydrogen available to carry off the carbon’s attachment energy will increase the sticking probability of carbon
as well as reducing the carbon photodissociation rate. Frequent collisions with hydrogen atoms will replenish the hydrogen layer, and will increase the rate of carbon accretion as well as protecting the nascent cluster from photodissociation. The rate of growth of a carbon cluster will then be limited by the rate at which it can capture the more abundant hydrogen. Further calculations are needed to quantify this effect.

On the other hand, if nucleation is impossible around an erupting nova, we must explain why dust seems to form when the ejected material has had time to reach the condensation radius. It may be that pre-existing nucleation sites lie around the nova before the eruption, too small to obstruct much light unless they grow (Malakpur 1977, Bode and Evans 1980, Stickland, et al. 1981, Albinson and Evans 1987; but see objections by Bode 1982). When the eruption begins, the nuclei inside the condensation radius evaporate, while those outside may survive the initial encounter with ejecta and begin to grow. Alternatively, nuclei may form in the dense gas of an accretion (or exterior excretion) disc, where they are shielded from direct irradiation by the nova. They may then be swept into the outward gas flow. Hypotheses of this type depend sensitively on details of mixing in complex fluid flows.

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Appendix: A Stochastic Approach to Cluster Evaporation

In the body of this paper we calculated the photodissociation rate of a small cluster assuming it absorbs only one photon. When the photon flux is high this assumption is not justified, for the cluster has a significant chance of absorbing a second photon before it loses all the energy of the first. This effect is unimportant for our grain growth calculations, for under such high photon fluxes photodissociation destroys all small particles, and no larger ones form. It may, however, be of interest for the evolution of larger grains in the presence of large photon fluxes and for the growth of pre-existing nuclei. It also provides a mechanism by which photons whose energy is individually less than the photodissociation threshold may cause photodissociation. In this Appendix all photons are assumed to have the same energy of 4 eV, a fair approximation for grains heated by an 8000°K black body nova spectrum, but the dissociation energy is still taken to be 6 eV.

The heating of a cluster by visible or ultraviolet radiation can be treated as a shot noise process, with the absorption of photons being a Poisson process and their cooling by infrared emission determining the response function. The cooling rate is directly proportional to the energy content of the cluster, so that its response is linear. The response function is \( E_p \exp(-bt) \), where \( E_p \) is the energy of the absorbed photon and \( b = 100 \, \text{sec}^{-1} \), appropriate to infrared emission by molecular vibrations. We need to determine the probability distribution of the internal energy of the cluster, given an average photon absorption rate \( \lambda \) and a cooling time \( 1/b \).

The solution to this statistical problem was given by Papoulis (1966), under the approximation that the cooling after a photon absorption follows an exponential law for a finite time \( T \) and that the energy then abruptly drops to zero. This approximation becomes exact in the limit \( T \to \infty \). The probability density \( f(E) \) is

\[
f(E) = \exp(-\lambda T) \sum_{k=0}^{\infty} g_k(E) \frac{(\lambda T)^k}{k!},
\]  

(A.1)

where \( g_k(E) \) is the probability density of the cluster’s energy, given that it has absorbed
$k$ photons in the time interval $T$.

The function $g_0(E) = \delta(E)$ and the other $g_k(E)$ can be derived in the following way: If a photon of energy $E_p$ was absorbed in the preceding interval $T$ its remaining contribution to the cluster will be an energy between $E_p$ and $E_p \exp(-bT)$; if it was absorbed a time $t$ ago its remaining contribution is $E_p \exp(-bt)$. For any $0 < t < T$ the probability $P(E)$ that such a photon contributes less than $E = E_p \exp(-bt)$ is $(T - t)/T = 1 + \ln(E/E_p)/(bT)$. The probability density $g_1(E)$ is

$$g_1(E) = \frac{dP}{dE} = \frac{1}{bTE}.$$  \hfill (A.2)

The probability density $g_2(E)$ is the convolution of $g_1(E)$ with itself and, in general,

$$g_k(E) = g_{k-1}(E) * g_1(E).$$  \hfill (A.3)

These $g_k$ can be calculated by Fourier transforming $g_1$, raising it to the $k$-th power, and inverting the transform.

In practice we consider only the first 100 terms in the infinite sum (A.1), which are sufficient if $(\lambda T)^{100} \ll 100!$, or $\lambda T < 100/e \approx 37$. For $T$, we chose a value equal to six half-lives, during which the energy decays by a factor of 64; i.e., $T = \ln(64)/b$. The introduction of $T$ is only a calculational device with no physical significance. Since the absorbed photons each have an energy of 4 eV and the emitted infrared photons energies of 0.1–0.2 eV, the discreteness of the infrared emission, which we ignore, imposes a physical cutoff at $t = O(T)$. Our $T$ merely represents the time it takes for a cluster with 4 eV of internal energy to radiate its last vibrational quantum. Because $T$ has no physical significance, we will present our results in terms of the physically meaningful term $\lambda \ln 2/b = \lambda T_{1/2}$, the number of photons absorbed in a cooling half life.

For large fluxes of photons, the probability density $f(E)$ should approach a Gaussian. The average cluster energy (in units of 4 eV, the energy of our photons) is $\lambda/b$. Its variance is $\lambda/(2b)$ (Papoulis 1966). Thus we can construct a Gaussian probability distribution with
\[ \sigma = \left( \frac{\lambda}{2b} \right)^{1/2}: \]
\[ g(E) = \frac{1}{(2\pi\sigma^2)^{1/2}} \exp\left( -\frac{(E - \langle E \rangle)^2}{2\sigma^2} \right). \quad (A.4) \]

Given \( g(E) \) we can calculate the mean evaporation rate from the evaporation rate as a function of energy given by the RRK theory of unimolecular reaction rates (5). In this appendix we set \( S = 3N - 5 \), which is correct for the linear dimers and trimers; for nonlinear clusters \( S = 3N - 6 \). However, the difference is small for the larger clusters; we also calculated evaporation rates for quadrimers, assuming \( S = 3N - 6 \), and found evaporation rates increased at most by a factor of 2.5. For clusters of 20 the difference was less than 30%. The uncertainties in the absorption cross-section, or whether very small clusters can absorb at all at energies of 4 eV, are of greater importance.

We calculate the mean evaporation rate of clusters containing \( N \) atoms and \( S \) modes in a given radiation field using (A.1) and the RRK theory dissociation rate (5). It is
\[ R = \int_{0}^{\infty} f(E)\tau^{-1}_d(E) dE. \quad (A.5) \]

The radiation intensity enters through the parameter \( \lambda \) in (A.1). This may be used in the kinetic equations for grain growth.

At low values of \( \lambda/b \), fluctuations are the dominant effect controlling the evaporation rate. If the mean energy is used for the calculation, clusters whose mean energy is less than the dissociation energy have an evaporation rate of zero. The true evaporation rate can be quite high in small clusters, because of fluctuations which carry the energy over the dissociation threshold. At higher fluxes and for larger clusters (which have a higher absorption rate), the Gaussian approximation becomes an accurate approximation to the true evaporation rate. At still higher flux levels, fluctuations become insignificant; the evaporation rates given by the mean energy, the Gaussian approximation to the energy distribution, and the true energy distribution all agree. Figure A.1 displays the regimes in \( (\lambda T_{1/2}, N) \) space in which the various approximations are valid. At the highest \( \lambda T_{1/2} \) it is adequate to treat the clusters as having their mean energy. At lower \( \lambda T_{1/2} \) the Gaussian
approximation (A.4) to the energy distribution is adequate, while for the smallest \( \lambda T_{1/2} \) it is necessary to use the complete stochastic \( f(E) \) as in (A.1).

We have calculated evaporation rates for clusters in the size range 2–640 carbon atoms (larger clusters would have high absorption rates and fluctuations would be unimportant) at distances \( r \) in the range \( 7.5 \times 10^{13} – 6 \times 10^{14} \) cm from a \( L = 10^{38} \) ergs sec\(^{-1} \) source of 4 eV monochromatic radiation. These parameters describe nova shells at the time when dust is observed to form. We assume an absorption cross section of \( 10^{-18} \) cm\(^2\) per atom. Given these parameters the flux at the dust-forming distance is about \( 10^{19} \) photons cm\(^{-2}\) sec\(^{-1}\). Dimers will absorb about 20 photons per second, while an 10-mer will absorb 100 photons per second. The clusters must begin forming before the dust, and hence possibly at somewhat smaller \( r \) and higher fluxes. We are therefore interested in a range of values of \( \lambda \) between 10 and \( 10^3 \) sec\(^{-1}\), and the time between photon absorptions is \( 10^{-3} – 10^{-1} \) sec.

Some sample results are shown in Table A.1 for \( N = 3 \) and \( N = 20 \). For \( \lambda T_{1/2} \leq 1 \) the average energy of a cluster is less than 6 eV, the dissociation energy. The evaporation rate calculated using the mean energy is zero, while the evaporation rates calculated for \( N = 3 \) in the Gaussian approximation and exactly are up to \( 10^{12} \) sec\(^{-1}\), because the fluctuations take some grains above the energy needed for dissociation. These extraordinarily high rates are unphysical, because they violate our tacit assumption that dissociation is rare enough that the energy distribution of a cluster is determined only by radiative processes. They imply actual dissociation rates approaching \( \lambda \), at which this assumption breaks down. Such dissociation rates exceed the rate of accretion of carbon atoms and would prevent any cluster growth. Our assumptions are valid and our calculational technique applies when dissociation rates are comparable to association rates, the circumstances in which quantitative results are needed.

The exact calculation gives a much higher value for the evaporation rate than the Gaussian approximation at low flux levels; the Gaussian approximation underestimates
the extent of the fluctuations. As the flux increases the Gaussian approximation and the exact method become closer; in fact they cross and the Gaussian approximation becomes slightly larger than the exact approach. For still higher fluxes the calculation using the mean energy approaches the other two values.

A similar pattern is followed by clusters of larger size (Table A.2 shows results for 20-mers). Here the Gaussian approximation and the exact calculation come to agree fairly well. Still higher fluxes are needed to bring the calculation using the mean energy into agreement with the others.
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| Atoms in Cluster | Sticking Coefficient |
|------------------|----------------------|
| 1                | 2.2 \times 10^{-7}   |
| 2                | 3.6 \times 10^{-6}   |
| 3                | 8.2 \times 10^{-5}   |
| 4                | 8.5 \times 10^{-4}   |
| 5                | 5.3 \times 10^{-3}   |
| 6                | 2.4 \times 10^{-2}   |
| 7                | 7.8 \times 10^{-2}   |
| 8                | 0.19                 |
| 9                | 0.38                 |
| 10               | 0.57                 |

Table 1: Sticking Coefficients
| Atoms in Cluster | Vibrational Degrees of Freedom | Dissociation Time (sec) |
|-----------------|-------------------------------|------------------------|
| 3               | 3                             | $1.1 \times 10^{-11}$  |
| 4               | 6                             | $7.2 \times 10^{-9}$    |
| 5               | 9                             | $1.1 \times 10^{-6}$    |
| 6               | 12                            | $7.0 \times 10^{-5}$    |
| 7               | 15                            | $2.6 \times 10^{-3}$    |
| 8               | 18                            | $6.1 \times 10^{-2}$    |
| 9               | 21                            | 1.0                    |
| 10              | 24                            | 14.                    |

Table 2: Dissociation Times from RRK Theory for $E = E_d$
| $\lambda T_{1/2}$ | Mean Energy | Gaussian Distribution | Exact Distribution |
|-----------------|-------------|-----------------------|-------------------|
| 0.167           | 0.0         | $1.4 \times 10^7$     | $4.6 \times 10^9$ |
| 0.333           | 0.0         | $5.7 \times 10^7$     | $3.8 \times 10^{10}$ |
| 0.500           | 0.0         | $6.0 \times 10^{10}$  | $1.3 \times 10^{11}$ |
| 0.667           | 0.0         | $2.3 \times 10^{11}$  | $3.2 \times 10^{11}$ |
| 0.833           | 0.0         | $5.4 \times 10^{11}$  | $6.3 \times 10^{11}$ |
| 1.000           | 0.0         | $1.0 \times 10^{12}$  | $1.1 \times 10^{12}$ |
| 1.167           | $1.7 \times 10^{11}$ | $1.6 \times 10^{12}$ | $1.6 \times 10^{12}$ |
| 1.333           | $7.5 \times 10^{11}$ | $2.4 \times 10^{12}$ | $2.3 \times 10^{12}$ |
| 1.500           | $1.7 \times 10^{12}$ | $3.2 \times 10^{12}$ | $3.1 \times 10^{12}$ |
| 1.667           | $2.9 \times 10^{12}$ | $4.2 \times 10^{12}$ | $4.0 \times 10^{12}$ |
| 1.833           | $4.2 \times 10^{12}$ | $5.2 \times 10^{12}$ | $4.9 \times 10^{12}$ |
| 2.000           | $5.5 \times 10^{12}$ | $6.2 \times 10^{12}$ | $5.9 \times 10^{12}$ |
| 2.167           | $6.8 \times 10^{12}$ | $7.2 \times 10^{12}$ | $6.9 \times 10^{12}$ |
| 2.333           | $8.2 \times 10^{12}$ | $8.3 \times 10^{12}$ | $8.0 \times 10^{12}$ |
| 2.500           | $9.5 \times 10^{12}$ | $9.3 \times 10^{12}$ | $9.1 \times 10^{12}$ |
| 2.667           | $1.1 \times 10^{13}$ | $1.0 \times 10^{13}$ | $1.0 \times 10^{13}$ |
| 2.833           | $1.2 \times 10^{13}$ | $1.1 \times 10^{13}$ | $1.1 \times 10^{13}$ |
| 3.000           | $1.3 \times 10^{13}$ | $1.2 \times 10^{13}$ | $1.2 \times 10^{13}$ |
| 3.167           | $1.4 \times 10^{13}$ | $1.3 \times 10^{13}$ | $1.3 \times 10^{13}$ |
| 3.333           | $1.5 \times 10^{13}$ | $1.4 \times 10^{13}$ | $1.4 \times 10^{13}$ |

Table A.1: Trimer Photodissociation Rates (sec$^{-1}$)
| $\lambda T_{1/2}$ | Mean Energy | Gaussian Distribution | Exact Distribution |
|--------------------|-------------|----------------------|-------------------|
| 0.167              | 0.0         | $1.6 \times 10^{-8}$ | 1.2               |
| 0.333              | 0.0         | $2.6 \times 10^{-3}$ | 71.               |
| 0.500              | 0.0         | 0.86                 | $9.2 \times 10^2$ |
| 0.667              | 0.0         | 35.                  | $6.1 \times 10^3$ |
| 0.833              | 0.0         | $5.1 \times 10^2$   | $2.8 \times 10^4$ |
| 1.000              | 0.0         | $4.1 \times 10^3$   | $9.7 \times 10^4$ |
| 1.167              | $1.2 \times 10^{-6}$ | $2.2 \times 10^4$ | $2.9 \times 10^5$ |
| 1.333              | $1.0 \times 10^{-3}$ | $8.8 \times 10^4$ | $7.3 \times 10^5$ |
| 1.500              | 0.11        | $2.9 \times 10^5$   | $1.7 \times 10^6$ |
| 1.667              | 3.6         | $8.1 \times 10^5$   | $3.6 \times 10^6$ |
| 1.833              | 58.         | $3.0 \times 10^6$   | $7.1 \times 10^6$ |
| 2.000              | $5.5 \times 10^2$ | $4.5 \times 10^6$ | $1.3 \times 10^7$ |
| 2.167              | $3.7 \times 10^3$ | $9.2 \times 10^6$ | $2.4 \times 10^7$ |
| 2.333              | $1.8 \times 10^4$ | $1.8 \times 10^7$ | $4.0 \times 10^7$ |
| 2.500              | $7.4 \times 10^4$ | $3.2 \times 10^7$ | $6.5 \times 10^7$ |
| 2.667              | $2.5 \times 10^5$ | $5.5 \times 10^7$ | $1.0 \times 10^8$ |
| 2.833              | $7.3 \times 10^5$ | $9.1 \times 10^7$ | $1.6 \times 10^8$ |
| 3.000              | $1.9 \times 10^6$ | $1.5 \times 10^8$ | $2.3 \times 10^8$ |
| 3.167              | $4.4 \times 10^6$ | $2.2 \times 10^8$ | $3.4 \times 10^8$ |
| 3.333              | $9.6 \times 10^6$ | $3.3 \times 10^8$ | $4.8 \times 10^8$ |
| 4.167              | $1.8 \times 10^8$ | $1.7 \times 10^9$ | $2.1 \times 10^9$ |
| 5.000              | $1.3 \times 10^9$ | $6.1 \times 10^9$ | $6.8 \times 10^9$ |

Table A.2: 20-mer Photodissociation Rates (sec$^{-1}$)
Figure Captions

Figure 1: Size distribution of grains found without photodissociation. $M_C$ is the carbon mass and $f$ the number of atoms in a grain. The units of the ordinate are arbitrary, but the relative calibration of the two cases (labeled by the total mass of carbon in the nova debris) is meaningful, as is the grain radius $a$.

Figure A.1: Regions of validity (10% accuracy) of exact, Gaussian, and mean energy calculations of photodissociation rates.