RADIO EMITTING DUST IN THE
FREE ELECTRON LAYER OF SPIRAL GALAXIES:
TESTING THE DISK/HALO INTERFACE

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ABSTRACT

We present a study of the radio emission from rotating, charged dust grains immersed in the ionized gas constituting the thick, Hα-emitting disk of many spiral galaxies. Using up-to-date optical constants, the charge on the grains exposed to the diffuse galactic UV flux has been calculated. An analytical approximation for the grain charge has been derived, which is then used to obtain the grain rotation frequency. Grains are found to have substantial radio emission peaked at a cutoff frequency in the range 10-100 GHz, depending on the grain size distribution and on the efficiency of the radiative damping of the grain rotation. The dust radio emission is compared to the free-free emission from the ionized gas component; some constraints on the magnetic field strength in the observed dusty filaments are also discussed. The model can be used to test the disk-halo interface environment in spiral galaxies, to determine the amount and size distribution of dust in their ionized component, and to investigate the rotation mechanisms for the dust. Numerical estimates are given for experimental purposes.
1. MOTIVATION

Interstellar dust is known to be responsible for redistributing the energy in the interstellar radiation field mainly by reemitting in the far-infrared (e.g., Mathis 1990). In the following it is demonstrated that the classical problem (Chandrasekhar 1943, Erickson 1957, Hoyle & Wickramasinghe 1970) of charged and spinning dust particles caused by interaction of grains with the surrounding gas and radiation field will lead to radiation at easily accessible radio continuum wavelengths if the distributions of grain sizes and charges are properly taken into account. As a consequence, it is well possible that in specific astrophysical environments this dust radio emission contributes to the observed radio continuum spectra in addition to thermal and non-thermal processes.

Of particular interest in this respect is the thick gas layer of spiral galaxies. The large scale-height of the, most probably, photoionized gas layer in some galaxies such as the Milky Way (Reynolds 1991) is indicative of an intensive UV-radiation field high above the plane. In addition, the presence of dust in these disk/halo interfaces of spiral galaxies on kpc scales has been suggested by theoretical (Franco et al. 1991) and observational work (Keppel et al. 1991). This is an interesting situation to consider the scenario developed in this paper. The application to such an environment could result in independent determinations of dust masses, grain size distributions, gas-to-dust ratios, etc.

2. GRAIN CHARGE

In this Section we will derive the charge, $Z$, on grains located in the Hα emitting gas known to extend to heights well in excess of the main gaseous disk of the parent spiral
galaxy. Grains interact with the surrounding (partially ionized) gas, with the radiation field photons and with other grains; we will neglect however the last possibility given the low collision rate of the process. Throughout the paper we will adopt the values given by Reynolds (1993) \( T_e = 10^4 \) K, \( n_e = 0.1 \text{ cm}^{-3} \) for the electron temperature and density in our Galaxy, respectively.

The UV radiation field at high galactic latitudes, the one relevant for the present calculation, has been subject of several studies in the last years and it is still subject of debate. A rather complete collection of observational results has been compiled by Reynolds (1990). The data are well fitted by a power-law extending approximately from 8 to 300 eV, with the flux expressed as

\[
F_\nu = F_{\nu_L} \epsilon_{Ry}^{-4},
\]

where \( F_{\nu_L} = 10^4 \text{ phot cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ eV}^{-1} \) is the flux at the Lyman limit and \( \epsilon_{Ry} \) is the photon energy expressed in \( Ry \); eq. (2.1) is also consistent with data not included in the collection (F. Paresce, private communication). Some uncertainty is present in the determination of the lower and upper limits to \( F_\nu \), but for our aims this does not represent a problem as the lower limit is dictated by the threshold energy for the photoelectric effect \( (B \sim 8 \text{ eV for most grain materials}) \) and the upper limit, \( \epsilon_{max} \gg B \), by the optical properties of the grains (above 300 eV they are virtually transparent). An useful parameter to be defined is \( \Xi = F_{\nu_L} / n_e T_e \) proportional to the ionization parameter.

The second necessary ingredient is the absorption coefficient, \( Q_{abs}(a, \epsilon) \), where \( a \) is the grain size, of the dust in the UV/soft X-rays range. We have calculated numerically \( Q_{abs}(a, \epsilon) \) through the Mie formulae for spherical silicate grains using the dielectric con-
stants kindly provided by Martin & Rouleau (unpublished, but some results are in Martin & Rouleau 1991). These constants are an extension of the Draine & Lee (1984) ones for the astronomical silicate up to 16 eV; from 16 eV to 31 eV they are extrapolated from Huffman (1975) and above 31 eV a synthesized FeMgSiO₄ cross section has been used. The results are shown in Fig. 1 for different grain sizes. In the range of the scattering parameter $x = 2\pi a/\lambda$ we are interested in, none of the classical physical approximations (Rayleigh-Gans, geometric optics) can be used for all the sizes. Since it is useful to have an analytical approximation anyway, we have fitted the numerical results with the following law:

$$Q_{\text{abs}}(a, \epsilon) = \left( \frac{\epsilon}{\delta \text{ eV}} \right)^{2.5[(a/0.1\mu m)^{0.1}-1]} = \left( \frac{\epsilon}{\delta \text{ eV}} \right)^{\gamma_a}$$ (2.2)

We will make use of this simple law, which is rather rough for large grains but it is reasonably accurate for $a \lesssim 0.1\mu m$, to calculate the average absorption coefficient weighted on the spectrum: $\langle Q_{\text{abs}}(a) \rangle = 3/(3 - \gamma_a)$.

We have solved the detailed balance equation for the grain charge numerically; the results are shown in Fig. 2 for different values of $F_L$. The explicit form of such an equation is given in the Appendix; here we just mention the processes that we have included in the calculation. In a pure hydrogen plasma, as the one we have assumed, positive charge is accumulated on the grain surface by i) proton collisions; ii) secondary electron emission following ion collisions; iii) photoelectric effect. Negative charging processes are i) electron collisions ii) secondary electron emission following electron collisions. Field emission, which limits the amount of charge on the grain, has been taken into account even if, in the conditions of interest, this limiting value has never been reached. Field ionization of the H
atoms coming from impinging protons that have recombined with \( \pi \)-electrons of the lattice is negligible for the grain electrostatic potential of interest here. Since thermal sputtering yields are negligible below \( T \lesssim 10^5 \) K, there are two mechanisms that may modify and limit the grain size distribution. Shattering due to electrostatic stress \( S_e \) could destroy very small grains: however, for the adopted parameters (\( \Xi = 10 \)) and a tensile stress \( \tau = 10^{10} \text{ ergs cm}^{-3} \) (Draine & Salpeter 1979) this does not affect grains with

\[
a \gtrsim 2 \times 10^{-8} \left( \frac{\tau}{10^{10} \text{ ergs cm}^{-3}} \right)^{-1/2} \left( \frac{\eta}{1.49 \times 10^7 \text{ ergs}} \right) \text{ cm}
\]

(\( \eta \) is defined in eq. [2.4]). The second possible effect is related to the centrifugal stress \( S_c \) arising from the grain rotation: we will discuss this point in the next Section.

The left panel of Fig. 2 shows the grain charge as a function of \( a \) for two different values of \( \mathcal{F}_{\nu L} \); in the right panel the rates for the various charging processes are shown for the standard case \( \mathcal{F}_{\nu L} = 10^4 \) (\( \Xi = 10 \)). Clearly, the stronger the radiation field the smaller is the value of the negative charge.

Looking at the various rates it is immediately evident that a simple approximation can be attempted: in this case, only two processes, \( i.e. \) ion and electron collisions, substantially contribute to the charge equilibrium. Thus we can find an approximate solution of the detailed balance (eq. [A1]) obtained considering only the two above mentioned processes. A simple linear relation between \( Z \) and \( a \) approximates very accurately the numerical solutions found (see Appendix):

\[
Z(a) = -2.504 \frac{k T_e}{e^2} a = -1.49 \times 10^7 \left( \frac{T_e}{10^4 \text{ K}} \right) a = -\eta a;
\]

for the adopted \( T_e, \eta = 1.49 \times 10^7 \) (cgs units). Eq. (2.3) is also a restatement of the
classical result obtained by Spitzer (1941) and it is indicative of the scarce relevance of the radiation field in determining the grain charge.

3. DUST RADIO EMISSION

Grains interact with the gas and the radiation field photons and they are put into rotation. There are a number of random torques that need to be considered in order to calculate the steady rms angular velocity $\omega$ of grains: collisions with gas atoms, absorbed or emitted photons, photoelectric emission and molecule formation on the grain surface. The various impulsive torques are supposed to act randomly in direction, thus the angular momentum grows in a random walk; if surface irregularities are present torques may be not time-averaged and therefore lead to a suprathermal rotation. In order to evaluate $\omega$ we will rely on the results obtained by Purcell (1979).

The rms rotational velocity $\langle \omega \rangle$ can be obtained in a straightforward manner once the effective temperature $T_{\text{eff}}$ characterizing the Brownian motion is known from

$$\frac{1}{2} I \omega^2 = \frac{3}{2} k T_{\text{eff}}, \quad (3.1)$$

where $I$ is the rotational inertia. If the only processes exciting rotation would be atom-grain collisions and atom evaporation from the grain surface then $T_{\text{eff}}$ would be exactly the arithmetic mean of $T_e$ and $T_d$, where $T_d$ is the dust temperature. Non-thermal contributions come from photoelectric emission, photon absorption/emission and $H_2$ molecule formation. Let $R$ be the ratio of the rates of increment of the grain mean square angular momentum about the rotation axis $z$, $\langle \Delta J_z^2 \rangle$, due to this process with respect to atom
collisions; then
\[ R = \frac{N_e \langle \Delta J_z^2 \rangle_e}{N_H \langle \Delta J_z^2 \rangle_H} = \frac{m_e}{m_H} \frac{N_e \langle \epsilon \rangle_e}{N_H \langle \epsilon \rangle_H}. \] (3.2)

In the previous formula \( N_e \) and \( N_H \) are the rates of photoelectric ejection and of atom collisions per unit area, respectively; \( \langle \epsilon \rangle_e \) and \( \langle \epsilon \rangle_H \) are the average kinetic energies corresponding to the two processes. For negative grain charge, \( N_e = 17.5 \mathcal{F}_{\nu,\infty} Q_{abs}(a) \sim 8.75 \times 10^4 \) (cgs units); \( \langle \epsilon \rangle \sim (\langle \epsilon \rangle - B)/2 \sim 2 \text{ eV} \), where \( \langle \epsilon \rangle \) is the spectrum average photon energy (Draine 1978); it follows \( R \sim 5 \times 10^{-3} \). We have checked the importance of rotational excitation due to diffuse galactic UV photons and \( H_2 \) molecule formation. The first one is completely negligible, while the second may very marginally contribute to \( T_{eff} \), since \( R_{H_2} \), analogous to the photoelectric \( R \), is \( R_{H_2} = 0.025 \). This low value can be understood observing that the average energy release following a \( H_2 \) molecule formation \( (\sim 0.2 \text{ eV}) \) is substantially smaller than \( kT_e \). Thus
\[ T_{eff} = \frac{1}{2} T_e \left[ 1 + \left( \frac{T_d}{T_e} \right) + \sum_j R_j \right] \approx \frac{1}{2} T_e, \] (3.3)

since \( T_d \ll T_e \); \( j \) runs through all the above mentioned processes.

Assuming spherical grains, upon substitution of eq. (3.3) into eq. (3.1), we obtain a simple relation between the radius and the rotational velocity
\[ a = 5.94 \times 10^{-4} T_{eff}^{1/5} \omega^{-\gamma} = 3.26 \times 10^{-3} \left( \frac{T_e}{10^4 \text{K}} \right)^{1/5} = \beta \omega^{-\gamma}, \text{ (cgs units)} \]

having assumed a grain material density \( \delta = 3.3 \text{gcm}^{-3} \) to eqs. (3.2)-(3.4) we will use the value \( \beta = 3.26 \times 10^{-3} \) (cgs units).

Finally, since grains are charged and rotating, radiative damping due to dipole emission may slow down the rotation rate. For a grain with an electric dipole moment \( d \simeq \xi Z e a \)
(\(\xi \sim 0.01\)) the radiative damping time becomes equal to the collisional damping time for

\[
a_d = 10^{-7} \left( \frac{\xi U}{V} \right)^{2/5} \left( \frac{n_e}{0.1 \text{cm}^{-3}} \right)^{-1/5} \left( \frac{\delta}{3.3 \text{g cm}^{-3}} \right)^{-1/5} \left( \frac{T_e}{10^4 \text{K}} \right)^{1/10} \text{cm},
\]

Hence, rotation of very small grains \((a \lesssim 10 \, \text{\AA})\) will be affected by radiative damping.

It is appropriate at this point to pose the question about the possible grain centrifugal disruption as well. Grains are destroyed if the centrifugal stress \(S_c \sim 3 I \omega^2 / 8 \pi a^3 > \tau\): thus, it must be \(\omega < \sqrt{\delta / \tau} \, a^{-1} = 1.23 \times 10^5 a^{-1} \text{ Hz}\). From a comparison with eq. (3.4), it appears that grains may easily survive the centrifugal stress.

There are several ways in which a grain may acquire dipole moments: for instance, if the grain is not perfectly spherical, it is rather possible that its center of mass and its center of charge do not coincide. However, even for a spherical grain a net electric dipole moment \(d\) is likely to arise from statistical fluctuations in the distribution of the charge trapped in or on the grain. A spherical grain with electric potential \(U = Z \varepsilon / a\) and capacity \(C = a\), has a number of charges on each hemisphere equal to \(Z_{1/2} = UC / 2 \varepsilon = Ua / 2 \varepsilon\); the statistical fluctuation of this number is roughly \(\sqrt{Ua / 2 \varepsilon}\). Thus, since the length of the dipole is \(\simeq a\), the average dipole moment is \(\langle d \rangle \simeq \sqrt{Z / 2 \varepsilon a}\). Note that, in this way, the electric dipole is only a few percent of the value that one would have obtained simply setting \(d = Z \varepsilon a\).

The volume emissivity \(\epsilon_v^d\) from the rotating, charged grains can be easily calculated once the dust size distribution is specified. We will assume for the latter a power law of the form

\[
\frac{dn}{da} = A_i n_H a^{-q} \quad \text{for} \quad a_0 \leq a \leq a_1,
\]

where \(n_H\) is the hydrogen density, which reduces to the standard Mathis et al. (1977) (MRN) distribution if \(q \sim 3.5\); in this case, the coefficient \(A_i = 3.3\pi \mu m_H \sigma / 8 \pi \delta \sqrt{a_1}\),
(where \( \mu \) is the mean molecular weight and \( \sigma \sim 0.01 \) is the dust-to-gas ratio) as obtained normalizing the distribution with \( \sigma \).

Using eqs. (2.3), (3.4), (3.6) we write

\[
\epsilon^d d\omega = \frac{1}{3} \frac{e^2 \eta}{c^3} \gamma \beta^{4-q} A_{\text{in\,H}} \omega^{3+\gamma+q} d\omega, \quad \text{for} \quad \left( \frac{a_1}{\beta} \right)^{\frac{1}{4}} \leq \omega \leq \left( \frac{a_0}{\beta} \right)^{-\frac{1}{4}}. \quad (3.7)
\]

Eq. (3.7) contains all the physics introduced in the calculation so far; its implications are discussed in the next §.

4. OBSERVATIONAL IMPLICATIONS

In the previous § we have predicted the amount of radio emission expected from dust immersed in a warm (\( T_c \sim 10^4 \) K), low-density (\( n_e \sim 0.1 \) cm\(^{-3} \)) gas, exposed to a diffuse UV spectrum. Here we evaluate the possibility to detect the predicted emission.

First we note that the wavelength range of the emission falls in a spectral band that can be easily observed. In fact, if we use \( a_1 = 0.25 \mu\text{m} \) and \( a_0 = 3 \) Å, as suggested from the previous analysis, emission is expected in the frequency range (\( \nu = \omega/2\pi \))

\[
3.1 \times 10^4 \leq \nu \leq 6.2 \times 10^{11} \quad \text{Hz}.
\]

In absence of radiative damping, small grains contribute to the 100 GHz band and the position of the high frequency cutoff \( \nu_c \) turns out to be a powerful tracer of the smallest sizes in the dust distribution; Fig. 3 (left panel) shows the dependence of \( \nu_c \) on \( a_0 \) for different temperatures. In reality the upper limit will be decreased to lower frequencies because of the radiative damping; in the present simple approach, however, we have not investigated the modification of the spectrum due to this effect. Note that \( \nu_c \) depends
very weakly on the model parameters (through $R$) and also on the temperature $T_e$. The knowledge of the minimum grain size (and possibly a test of the MRN distribution) would be a highly valuable piece of information on the dust properties.

There is no doubt that the detection of the radio emission we predict would be the major consequence of this paper; to evaluate this possibility we compare the emissivity per H nucleus given by eq. (3.7) with the free-free emissivity, $\epsilon_{\nu}^{ff}$ in the same frequency range by the ionized gas surrounding the dust. We suppose that the medium is optically thin, which is certainly a good approximation either for the gas and the dust at these frequencies; then

$$\epsilon_{\nu}^{d} = 1.36 \times 10^{-37} \left( \frac{\nu}{100 \text{ GHz}} \right)^{2.8} \text{ergs s}^{-1} \text{ Hz}^{-1};$$  

$$\epsilon_{\nu}^{ff} = 6.78 \times 10^{-41} \left( \frac{n_e}{0.1 \text{ cm}^{-3}} \right) \left( \frac{T_e}{10^4 \text{ K}} \right)^{-1/2} g(\nu, T_e) \text{ ergs s}^{-1} \text{ Hz}^{-1},$$  

where $g(\nu, T_e) = (\sqrt{3}/\pi) \log(5 \times 10^7 T^{3/2}/\nu)$. The comparison between the two fluxes is shown in the right panel of Fig. 3.

The dust spectrum should present a steep positive spectral index $\alpha$ which depends only on the size distribution index $q$, since $\gamma$ is fixed and determined by the microphysics; for a MRN distribution $\alpha = 2.8$. This should reduce the possibility of confusion with the flat free-free spectrum and Fig. 3 clearly demonstrates that the emission can in principle be detected.

An intriguing hypothesis can be suggested about the observed vertical dust filaments emerging from many spiral galaxy disks and immersed in the H$\alpha$ emitting gas (Sofue 1987; for a review see Dettmar 1992). We propose that these filamentary shapes arise because a magnetic field $B$ aligned with the filament axis is present and the charged dust
is tightly coupled to the field. Such a vertical configuration of the field could be produced by the enhancement of the Parker instability driven by radiation pressure on charged grains (Ferrara 1993). Thus, the field lines act as tracks for the grain motion and, vice versa, dust would represent an excellent tracer of the B-field topology. A necessary condition for this to be true is that the Larmor radius, $r_L$, for the grains must be much smaller than the transverse size of the filaments. Using the previous results, the maximum value of $r_L$ is

$$r_L = \frac{4\pi\delta e v}{3e\mu B} a_1^2 = 0.12 \left( \frac{v}{100\text{ km s}^{-1}} \right) \left( \frac{B}{\mu\text{G}} \right)^{-1} \left( \frac{a_1}{250\text{ nm}} \right)^2 \text{ pc.} \quad (4.3)$$

Thus, even for models predicting high velocities of the grains due to radiative acceleration (Ferrara et al. 1991), grains are strictly tied to the magnetic lines. A definite test of this interpretation would be represented by optical polarization measurements in which a correlation between the degree of polarization due to aligned grains and extinction toward background objects is found.

**APPENDIX**

**THE DETAILED BALANCE EQUATION**

The detailed balance equation we have used to determine the grain charge $Z$ is

$$n_p \left\{ \left( \frac{8kT_e}{\pi m_p} \right)^{1/2} \left[ J_{\nu_\tau_p}(a, Z, T_e) s_p + \delta_p(a, Z, T_e) \right] + J_{p_e}(a, F_\nu) \right\} =$$

$$n_e \left\{ \left( \frac{8kT_e}{\pi m_e} \right)^{1/2} \left[ J_{\nu_\tau_e}(a, Z, T_e) s_e - \delta_e(a, Z, T_e) \right] \right\}, \quad (A1)$$

where the index $p,e$ stands for protons and electrons, respectively; $J_{\nu_\tau}$ are the collisional charging rates; $\delta$ are the secondary emission rates; $J_{p_e}$ is the photoelectric charging rate; $s$ is the sticking probability. $J_{\nu_\tau}$ are taken from Draine & Sutin (1987); $J_{p_e}$ is taken from
Draine (1978); \( \delta \) and \( s \) are taken from Draine & Salpeter (1979). Eq. (A1) must be solved together with the field emission condition (Draine & Sutin 1987) that limits the value of the charge:

\[
-1 - 0.7 \left( \frac{a}{\text{nm}} \right)^2 < Z < 1 + 21 \left( \frac{a}{\text{nm}} \right)^2.
\] (A2)

When only ion and electron collisions are considered in a pure H plasma, eq. (A1) can be written as

\[
\left( \frac{m_e}{m_p} \right)^{1/2} = \frac{\tilde{J}_{\nu, \tau}(a, Z, T_e)}{\tilde{J}_{e, \tau}(a, Z, T_e)},
\] (A3)

with \( s_e = s_p = 1 \). For \( |\nu| \ll 1 \), \( \tilde{J} \to (1 - \nu/\tau) \) for \( \nu < 0 \) and \( \tilde{J} \to e^{-\nu/\tau} \) for \( \nu > 0 \), where \( \nu = Z e / q \), \( \tau = a k T / q^2 \) (see Draine & Sutin 1984) and \( q \) is the charge of the colliding particle. Thus

\[
\left( \frac{m_e}{m_p} \right)^{1/2} = \frac{e^\psi}{1 - \psi},
\] (A4)

where \( \psi = Z / \tau \). The solution of eq. (A4) is \( \psi = -2.504 \) and eq. (2.3) follows accordingly.

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REFERENCES

Chandrasekhar, S. 1943, Rev. Mod. Phys., 15, 1
Dettmar, R.-J. 1992, Fundamental of Cosmic Physics, 15, 143


Mathis, J. S. 1990, ARAA, 28, 37


Reynolds, R. J. 1990, The Galactic and Extragalactic Background Radiation, eds. S.
Bowyer & C. Leinert, (Dordrecht:Kluwer), 165


Sofue, Y. 1987, PASJ 39, 547

**FIGURE CAPTIONS**

**Figure 1** $Q_{ab}(a, e)$ for spherical silicate grains using the dielectric constants given by Martin & Rouleau; numbers indicate the grain radius in nm.

**Figure 2** Numerical solutions of the detailed balance equation (see Appendix). **Left:** grain charge as a function of $a$ for different values of $F_{\nu_L}$, as shown by numbers; **right:** rates for the various charging processes for the standard case $F_{\nu_L} = 10^4$; *long-dashed line* - electron collisions; *solid* - proton collisions; *short-dashed* - photoelectric effect; *dotted* - proton secondary emissions; *dot-dashed* - electron secondary emission.

**Figure 3** **Left:** dependence of the cutoff frequency $\nu_c$ on the minimum size of the grain distribution, $a_0$, for different temperatures: *solid* - $T_e = 10^4$ K; *dotted* - $T_e = 7000$ K; *dashed* - $T_e = 1.2 \times 10^4$ K. **Right:** comparison between the free-free emissivity (*flat dotted line*) from a gas with $T_e = 10^4$ K and $n_e = 0.1$ cm$^{-3}$ and dust emissivity (*steep dotted line*) per H nucleus; the *solid* line is the sum of the two. The dashed region denotes the range of frequency in which the spectrum may be modified by radiative damping (see text). In the upper right corner the cutoff frequency $\nu_c$ is shown.