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Radiation Pressure and the Poynting–Robertson Effect for Fluffy Dust Particles

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A correct understanding of the dynamical effect of solar radiation exerted on fluffy dust particles can be achieved with assistance of a light scattering theory as well as the equation of motion. We reformulate the equation of motion so that the radiation pressure and the Poynting–Robertson effect on fluffy grains are given in both radial and nonradial directions from the center of the Sun. This allows numerical estimates of these radiation forces on fluffy dust aggregates in the framework of the discrete dipole approximation, in which the first term of the scattering coefficients in Mie theory determines the polarizability of homogeneous spheres forming the aggregates.

The nonsphericity in shape turns out to play a key role in the dynamical evolution of dust particles, while its consequence depends on the rotation rate and axis of the grains. Unless a fluffy dust particle rapidly revolves on its randomly oriented axis, the nonradial radiation forces may prevent, apart from the orbital eccentricity and semimajor axis, the orbital inclination of the particle from being preserved in orbit around the Sun. However, a change in the inclination is most probably controlled by the Lorentz force as a consequence of the interaction between electric charges on the grains and the solar magnetic field. Although rapidly and randomly rotating grains spiral into the Sun under the Poynting–Robertson effect in spite of their shapes and structures, fluffy grains drift inward on time scales longer at submicrometer sizes and shorter at much larger sizes than spherical grains of the same sizes. Numerical calculations reveal that the dynamical lifetimes of fluffy particles are determined by the material composition of the grains rather than by their morphological structures and sizes. The Poynting-Robertson effect alone is nevertheless insufficient for giving a satisfactory estimate of lifetimes for fluffy dust grains since their large ratios of cross section to mass would reduce the lifetimes by enhancing the collisional probabilities. We also show that the radiation pressure on a dust particle varies with the orbital velocity of the particle but that this effect is negligibly small for dust grains in the Solar System.

Key Words: **interplanetary dust; meteoroids; orbits; solar radiation; zodiacal light.**

1. INTRODUCTION

Dust particles in the Solar System interact with the electromagnetic radiation field of the Sun as well as the solar gravitational field. When a dust particle is assumed to be spherical with a homogeneous structure, the solar radiation exerts a force on the particle in the direction of propagation of solar radiation, antiparallel to solar gravity. Both radiation pressure and gravity acting on a homogeneous sphere approximately obey the inverse square law of distances from the center of the Sun. The ratio β of radiation pressure to gravity on a particle is, therefore, a useful dimensionless quantity to evaluate the relative importance of radiation pressure in the dynamical evolution of dust particles. The values of β ratios for homogeneous spherical particles in the Solar System and extra-solar systems have been computed on the basis of Mie theory, which provides the exact solution of Maxwell's equations for a homogeneous sphere (Mukai and Mukai 1973, Artymowicz 1988, Lamy and Perrin 1997). There is a general agreement that the β ratio for a sphere reaches its maximum at submicrometer sizes in the Solar System, while the magnitude of the β ratio depends on the material of the particle (Burns *et al.* 1979). Clearly radiation pressure plays a key role in the dynamical behavior of submicrometer-size grains in the solar radiation and gravitational fields.

The radiation pressure depends not only on the size and material of dust grains but also on their shapes and structures (Gustafson 1989). As a natural consequence of coagulation growth, primordial grains in the solar nebula are assumed to have

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an aggregate structure (Weidenschilling *et al.* 1989). Because such aggregate particles are expected to have formed comets during the formation of the Solar System, dust grains released from comets are inevitably fluffy aggregates. Greenberg and Gustafson (1981) have proposed a model of cometary dust as agglomeration of elongated submicrometer-size interstellar dust. The KOSI-9 experiments simulating a cometary nucleus have shown that dust particles ejected from the surface of an ice–dust mixture are highly fluffy agglomeration of small individual grains (Grün *et al.* 1993, Thiel *et al.* 1995). Comets are also thought to be a source of interplanetary dust particles that have been observed as the zodiacal light. Giese *et al.* (1978) concluded from an analysis of zodiacal light brightness and polarization that interplanetary dust particles are fluffy. In fact, a typical structure of interplanetary dust particles collected in the Earth's upper stratosphere is a fluffy aggregate, which consists of many small grains ranging from less than 0.01 μm to several micrometers (Brownlee *et al.* 1980, Jessberger *et al.* 2001). Alternatively, interplanetary dust particles may originate from asteroids as a result of hypervelocity impacts of meteoroids onto regolith of the asteroids. Analogous to lunar glassy agglutinates, which are produced by impacts of micrometeoroids onto the lunar regolith, asteroidal grains could also be fluffy (Pillinger 1979, Keller *et al.* 1993). Levasseur-Regourd *et al.* (1997) claim that interplanetary, cometary, and asteroidal dust particles are essentially fluffy aggregates, because they found that the phase curves of linear polarization measured for the zodiacal light, cometary comae, and asteroidal regoliths are similar to computational results for highly fluffy aggregates. It is clear that a study of the radiation pressure on fluffy dust particles helps to better understand the dynamics of dust in the Solar System.

The radiation pressure on fluffy dust grains has been estimated by numerical computations based on microwave analogue measurements or based on Mie theory with a combination of effective medium approximations (EMAs), such as the Maxwell–Garnett mixing rule and the Bruggeman mixing rule (Mukai *et al.* 1992, Gustafson 1994, Wilck and Mann 1996, Kimura *et al.* 1997, Kimura and Mann 1999a, Gustafson *et al.* 2001). The application of EMAs relies on the assumption that a fluffy grain can be treated as a sphere having an average refractive index of constituent material and vacuum. As noted by van de Hulst (1957), not only does the radiation pressure on a particle have a component of force in the radial direction from the Sun, it also has a nonradial component, which never appears in the framework of Mie theory. The nonradial component of radiation pressure has been evaluated for infinite circular cylinders, spheroidal particles, and a plane mirror (Cohen and Alpert 1980, Voshchinnikov and Il'in 1983a,b, Voshchinnikov 1990, Klačka 1993, Il'in and Voshchinnikov 1998). In contrast to such specific cases of nonspherical particles, the discrete dipole approximation (DDA) originally proposed by Purcell and Pennypacker (1973) enables us to compute the radial and nonradial components of radiation pressure on fluffy dust particles (Kimura and Mann 1998, 1999b).

When a dust particle is moving with respect to the source of radiation, the equation of motion explicitly contains velocity-dependent terms, known as the Poynting–Robertson effect (Robertson 1937, Soter *et al.* 1977, Klačka 1992, Srikanth 1999). The Poynting–Robertson effect for a spherical dust grain causes a drag force along the opposite direction of dust motion as well as a Doppler effect in the direction of incident radiation (Robertson and Noonan 1968). These effects appear in the equation of motion to terms of order v/c, where v and c are the speed of the particle and that of light, respectively. As far as a spherical particle is concerned, the Poynting–Robertson effect gradually reduces the orbital eccentricity and semimajor axis due to the drag force, although the orbital inclination remains unchanged (Wyatt and Whipple 1950). Analogously to the radiation pressure acting on nonspherical dust particles, the direction of the Poynting–Robertson effect for fluffy dust particles may be different from that for homogeneous spherical grains. Nevertheless, no discussion has been included on the Poynting–Robertson effect for fluffy grains, because the evaluation of the Poynting–Robertson effect has been based on the equation of motion for spherical dust particles.

To correctly understand the radiation pressure and the Poynting–Robertson effect for fluffy dust particles, we first need to derive the equation of motion without assuming any specific shape of the particles. Second it is necessary to estimate quantitatively the radial and nonradial components of the radiation pressure forces on fluffy grains. Klačka (1994) actually attempted to generalize the equation of motion for arbitrarily shaped particles by introducing three unknown parameters concerning the interaction of dust particles with the solar radiation field. Unfortunately those parameters prevent us from calculating the radiation pressure and the Poynting–Robertson effect on dust particles of realistic composition, size, shape, and structure using available light-scattering theories.

The purpose of this study is to gain insight into the radiation pressure and the Poynting–Robertson effect on fluffy dust particles. Therefore we first describe the equation of motion for fluffy dust particles in radiation and weak gravitational fields so that the radiation pressure and the Poynting–Robertson effect for fluffy dust particles can be computed on the basis of a light-scattering theory. Next the radial and nonradial components of radiation pressure acting on fluffy aggregates of different material compositions are calculated using the DDA arranged for fluffy aggregates of spherical monomers. Finally we discuss the dynamical behavior of fluffy dust particles based on the numerical values for the radiation forces.

2. THE EFFECT OF RADIATION FORCES ON THE PARTICLE MOTION

2.1. Equation of Motion

In Appendix A, we describe the equation of motion for arbitrary-shaped particles by introducing the matrix representation for the radiation pressure cross section. An analogy between the equation of motion for arbitrary-shaped particles and that for spherical particles is clearly demonstrated, but it is not straightforward to compute the radiation pressure and the Poynting–Robertson effect on fluffy dust particles. We shall, therefore, reformulate the equation of motion, in which conventional notation for the asymmetry parameter and the extinction and scattering cross sections is used to describe the radiation forces. This formulation enables us to calculate easily the radiation forces using available light-scattering theories. Furthermore we provide the equation of motion as a superposition of orthogonal vectors so that the radial and nonradial components of the radiation forces can be discussed separately. Hereafter the various quantities involved in the equation of motion are taken to terms of order in v/c since we focus on the dynamics of dust in the Solar System. The complete relativistic equation of motion can be derived in the same manner but is beyond the scope of this paper. Although we deal with dust particles in the Solar System, our arguments can be easily extended to the cases for other stellar systems.

In the reference system of the Sun, we shall define orthogonal unit vectors $\mathbf{e}_1$, $\mathbf{e}_2$, and $\mathbf{e}_3$ (i.e., $\mathbf{e}_1 \cdot \mathbf{e}_2 = \mathbf{e}_2 \cdot \mathbf{e}_3 = \mathbf{e}_3 \cdot \mathbf{e}_1 = 0$). Let $\mathbf{e}_3 \cdot \mathbf{v} = 0$, where $\mathbf{v}$ denotes the velocity of a dust grain with respect to the Sun so that $\mathbf{e}_1$ and $\mathbf{e}_2$ determine the instantaneous orbital plane of the particle. Furthermore we set $\mathbf{e}_1 = \mathbf{e}_r$, where $\mathbf{e}_r$ is a unit vector parallel to the wave vector $\mathbf{k}$ of incident solar radiation in the reference system of the Sun. In the reference frame of the particle, we introduce orthogonal unit vectors $\mathbf{e}'_r$, $\mathbf{e}'_\theta$, and $\mathbf{e}'_\phi$ in a similar way. Namely, $\mathbf{e}'_r \cdot \mathbf{e}'_\theta = \mathbf{e}'_\theta \cdot \mathbf{e}'_\phi = \mathbf{e}'_\phi \cdot \mathbf{e}'_r = 0$, $\mathbf{e}'_\phi \cdot \mathbf{v} = 0$, and $\mathbf{e}'_r$ is parallel to the wave vector $\mathbf{k}'$ of incident solar radiation in the reference frame of the particle. The relation of the orthogonal unit vectors in those reference frames are given as follows:

$$\mathbf{e}'_r = -\frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_2 + \mathbf{e}_1, \tag{1}$$

$$\mathbf{e}'_\theta = \mathbf{e}_2 + \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_1, \tag{2}$$

$$\mathbf{e}'_\phi = \mathbf{e}_3. \tag{3}$$

Assuming that all the energy absorbed by a particle is isotropically reemitted at the same rate in the particle frame of reference, the particle in the solar frame of reference gains the energy E per unit time t as follows (see Eq. (A3); cf. Robertson and Noonan 1968):

$$\frac{dE}{dt} = \mathbf{v}\cdot\frac{d\mathbf{p}'}{dt}. \tag{4}$$

The momentum $\mathbf{p}'$ delivered to the particle per unit time in the reference frame of the particle and the delivered momentum $\mathbf{p}$ in the reference frame of the Sun fulfill the following equation (see Eq. (A4); cf. Robertson and Noonan 1968):

$$\frac{d\mathbf{p}}{dt} = \frac{d\mathbf{p}'}{dt}. \tag{5}$$

This quantity corresponds to the "time-averaged radiation pressure force" of Draine and Weingartner (1996) when the proper inertial reference frame of a particle is considered. Accordingly, the right-hand side of Eq. (5) is expressed as

$$\frac{d\mathbf{p}'}{dt} = U'(C'_{\rm ext}\mathbf{e}'_r - C'_{\rm sca}\mathbf{g}'), \tag{6}$$

where the energy density U' in the reference frame of the particle is given by (see Eq. (A6); cf. Robertson and Noonan 1968)

$$U' = U\left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right). \tag{7}$$

While $C'_{\rm ext}$ and $C'_{\rm sca}$ denote the usual extinction and scattering cross sections, the asymmetry parameter vector $\mathbf{g}'$ is defined by (Kimura and Mann 1998)

$$\mathbf{g}' = \frac{1}{C'_{\rm sca}}\int \mathbf{n}'\frac{dC'_{\rm sca}}{d\chi'}\,d\chi', \tag{8}$$

where $d\chi'$ is the element of solid angle, $\mathbf{n}'$ is a unit vector in the direction of scattering, and $dC'_{\rm sca}/d\chi'$ is the differential scattering cross section. Cartesian coordinates $(g'_r, g'_\theta, g'_\phi)$ describe the asymmetry parameter vector $\mathbf{g}'$ as

$$\mathbf{g}' = g'_r\mathbf{e}'_r + g'_\theta\mathbf{e}'_\theta + g'_\phi\mathbf{e}'_\phi, \tag{9}$$

where g'_r is the conventional asymmetry parameter. It is worthwhile noting that the nonradial components of radiation pressure force acting on fluffy dust grains originate from the $\mathbf{e}'_\theta$ and $\mathbf{e}'_\phi$ components of the asymmetry parameter vector.

From Eqs. (1)–(7) and (9), we obtain the equation of motion for a fluffy dust particle having mass m in the solar radiation field,

$$\begin{aligned} m\frac{d\mathbf{v}}{dt} = U\Bigg\{&\left[C'_{\rm pr}\left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right) - C'_{\rm sca}g'_\theta\frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\right]\mathbf{e}_1 \\ &- \left[C'_{\rm sca}g'_\theta\left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right) + C'_{\rm pr}\frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\right]\mathbf{e}_2 \\ &- C'_{\rm sca}g'_\phi\left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right)\mathbf{e}_3\Bigg\}, \end{aligned} \tag{10}$$

where $C'_{\rm pr} = C'_{\rm ext} - C'_{\rm sca}g'_r$ is the usual radiation pressure cross section (Draine 1988). If we regard $(C'_{\rm sca}g'_\theta, C'_{\rm sca}g'_\phi, C'_{\rm pr})$ as $A'(m, n, 1 - z)$, then Eq. (10) is equivalent to the equation of motion derived by Klačka (1994). In Appendix B, Eq. (10) is proved to be identical to Eq. (A11), which is the equation of motion with the matrix description of radiation pressure cross section.

2.2. Ratios of Radiation Forces to Gravity

To take into account the fact that the energy density U depends on a wavelength of the incident radiation, we integrate the right-hand side of Eq. (10) over the whole range of wavelengths. The energy density of the solar radiation at a wavelength range from λ to $\lambda + d\lambda$ is expressed as

$$dU = \pi \left(\frac{R_\odot}{r}\right)^2 \frac{B_\odot}{c}\, d\lambda, \tag{11}$$

where $R_\odot$ denotes the radius of the Sun and $B_\odot$ is the solar radiance at a wavelength of λ. The gravitational field of the Sun is so weak that the Newtonian approximation is applied (Landau and Lifshitz 1951). The gravitational force $\mathbf{F}_\mathrm{G}$ acting on a particle is therefore expressed to terms of order in v/c as

$$\mathbf{F}_\mathrm{G} = -\frac{GM_\odot m}{r^2}\mathbf{e}_r, \tag{12}$$

where G, $M_\odot$, and r are the gravitational constant, the mass of the Sun, and the distance of the particle from the center of the Sun, respectively. We finally can reformulate the equation of motion for fluffy dust particles in the radiation and gravitational fields of the Sun as

$$\begin{aligned}\frac{d\mathbf{v}}{dt} = &-\frac{GM_\odot}{r^2}\mathbf{e}_1 + \beta_1\frac{GM_\odot}{r^2}\left[\left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right)\mathbf{e}_1 - \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_2\right]\\ &+\beta_2\frac{GM_\odot}{r^2}\left[\frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_1 + \left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right)\mathbf{e}_2\right]\\ &+\beta_3\frac{GM_\odot}{r^2}\left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right)\mathbf{e}_3,\end{aligned} \tag{13}$$

where

$$\beta_1 = \frac{\pi {R_\odot}^2}{GM_\odot mc}\int_0^\infty B_\odot(C'_\mathrm{ext} - C'_\mathrm{sca}g'_r)\, d\lambda, \tag{14}$$

$$\beta_2 = \frac{\pi {R_\odot}^2}{GM_\odot mc}\int_0^\infty B_\odot(-C'_\mathrm{sca}g'_\theta)\, d\lambda, \tag{15}$$

and

$$\beta_3 = \frac{\pi {R_\odot}^2}{GM_\odot mc}\int_0^\infty B_\odot(-C'_\mathrm{sca}g'_\phi)\, d\lambda. \tag{16}$$

The β ratios describe the importance of the radiation forces on the dynamical evolution of dust particles relative to that of the solar gravitational force. We utilize the observational data compiled in Mukai (1990) for the solar radiance $B_\odot$ and Romberg's method for integration over wavelengths of the solar radiation, $\lambda = 0.14 - 300$ μm (Press *et al.* 1986). Because only the β ratios in Eq. (13) contain information on the properties of dust particles, we shall hereafter focus on numerical calculations of the β ratios alone based on the DDA, which is the most flexible method to derive numerically the cross sections for extinction C'_ext and scattering C'_sca, and the asymmetry parameter $\mathbf{g}'$ (Draine and Weingartner 1996).

3. DISCRETE DIPOLE APPROXIMATION

3.1. Formalism of Optical Properties

For the sake of simplicity, we shall omit the prime notation throughout Section 3, noting that the quantities described in the following are measured in the reference frame of a particle. In the framework of the DDA, fluffy dust particles are divided into N dipoles. Then $3N$ complex linear equations are solved to obtain a self-consistent set of dipole moment $\mathbf{P}_j$ for each dipole ($j = 1, \ldots, N$) so that $\mathbf{P}_j = \alpha_j\mathbf{E}_j$, where $\mathbf{E}_j$ is the electric field at the location $\mathbf{r}_j$ of the jth dipole and α_j is its polarizability. The extinction cross section C_ext and the scattering cross section C_sca are computed by (Draine 1988)

$$C_\mathrm{ext} = \frac{4\pi k}{|\mathbf{E}_0|^2}\sum_{j=1}^{N}\mathrm{Im}(\mathbf{E}^*_{\mathrm{inc},j}\cdot\mathbf{P}_j), \tag{17}$$

$$C_\mathrm{sca} = \frac{k^4}{|\mathbf{E}_0|^2}\int\left|\sum_{j=1}^{N}[\mathbf{P}_j - \mathbf{n}(\mathbf{n}\cdot\mathbf{P}_j)]e^{-ik\mathbf{n}\cdot\mathbf{r}_j}\right|^2 d\chi, \tag{18}$$

where $\mathbf{E}_{\mathrm{inc},j} = \mathbf{E}_0\exp(i\mathbf{k}\cdot\mathbf{r}_j - i\omega t)$ is the incident wave at $\mathbf{r}_j$. The asymmetry parameter vector $\mathbf{g}$ is calculated by (Draine 1988, Kimura and Mann 1998)

$$\mathbf{g} = \frac{\int\mathbf{n}\left|\sum_{j=1}^{N}[\mathbf{P}_j - \mathbf{n}(\mathbf{n}\cdot\mathbf{P}_j)]e^{-ik\mathbf{n}\cdot\mathbf{r}_j}\right|^2 d\chi}{\int\left|\sum_{j=1}^{N}[\mathbf{P}_j - \mathbf{n}(\mathbf{n}\cdot\mathbf{P}_j)]e^{-ik\mathbf{n}\cdot\mathbf{r}_j}\right|^2 d\chi}. \tag{19}$$

As mentioned by Draine (1988), the determination of the dipole polarizability α_j is an important input parameter for the DDA that is used to calculate the cross sections and the asymmetry parameter.

3.2. Determination of Dipole Polarizabilities

Draine and Goodman (1993) have shown that the lattice dispersion relation for the determination of dipole polarizabilities gives the most accurate results when a particle is divided into cubic cells. If a great number of dipoles on a lattice form into a fluffy dust particle, then there are huge computational memory and CPU time requirements (Draine and Flatau 1994). Okamoto (1995, 1996) has proposed the a1-term method to diminish the number of dipoles required for a cluster of spherical constituent particles. The a1-term method replaces the spherical monomers by single dipoles whose polarizability α_j is determined by

$$\alpha_j = i\frac{3a_1}{2k^3}, \tag{20}$$

where a_1 is the first scattering coefficient in Mie theory and is a function of the size parameter of the spherical monomer and the complex refractive index of the grain material (Bohren and Huffman 1983). Okamoto and Xu (1998) have shown that calculation results of extinction and scattering cross sections for small clusters of spherical monomers obtained by the a1-term method are in good agreement with the analytical solutions for the sphere clusters derived by Xu (1995). Therefore, we consider a cluster of spherical monomers as a representative model for fluffy dust particles to take advantage of the a1-term method.

4. MODEL FOR FLUFFY DUST PARTICLES

4.1. Morphology of Dust Particles

We specify the configuration of spherical monomers in clusters by adopting two cases of coagulation process: ballistic particle–cluster aggregation (BPCA) and the ballistic cluster–cluster aggregation (BCCA). The BCCA has been used to model dust grains in the primordial solar nebula as a consequence of cluster–cluster growth (Meakin and Donn 1988, Weidenschilling *et al.* 1989, Weidenschilling and Cuzzi 1993). The shape and structure of cometary dust particles simulated by KOSI-9 experiments were fairly similar to those of particles simulated by the BCCA (Grün *et al.* 1993). On the other hand, fluffy aggregates in interplanetary dust collections may have a resemblance to BPCA particles in shape and porosity (Brownlee *et al.* 1980). Accordingly, we expect that the BPCA and BCCA particles are equipped with the characteristics of dust particles in the Solar System.

Kitada *et al.* (1993) have performed computer simulations to produce the BPCA and BCCA particles in a three-dimensional space where the individual constituent spheres are assumed to be identical in size and to have the sticking probability of unity. The number N of spherical monomers and the radius of gyration s_g of an aggregate satisfy the relation $N \propto s_g^D$, which is known to be fractal. The fractal dimension D of the aggregates is found to be $D \approx 3$ for the BPCA and $D \approx 2$ for the BCCA. The particle density ρ of fractal aggregates can be described as $\rho \propto s_g^{D-3}$ (Smirnov 1990). Therefore the BCCA results in a highly porous and fluffy structure, compared to the BPCA (see Fig. 1).

4.2. Material Composition of Dust Particles

The extinction and scattering cross sections and asymmetry parameter of dust particles appearing in Eqs. (14)–(16) are a function of the material composition through the complex refractive indices. To investigate the dependence of material composition on the radiation pressure and the Poynting–Robertson effect, we take into account two types of material: silicate (weakly absorbing material) and carbon (strongly absorbing material). While carbon grains certainly exist in circumstellar envelopes

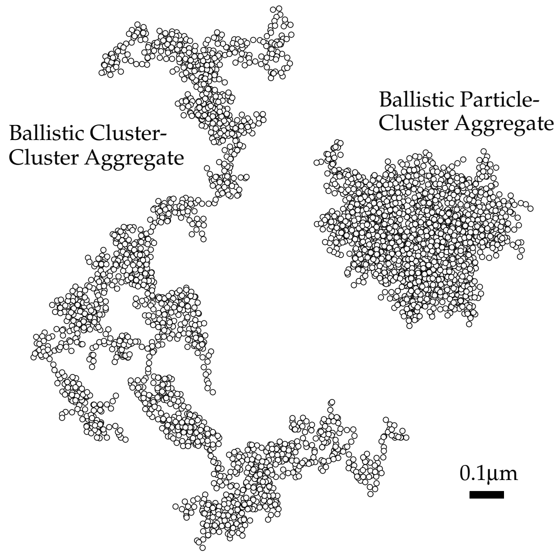

FIG. 1. An example of the ballistic particle–cluster aggregate and the ballistic cluster–cluster aggregate used in the calculations. The number N and radius s_{m} of the constituent particles are $N = 2048$ and $s_{\mathrm{m}} = 0.01$ μm, respectively.

of carbon stars, interplanetary dust particles are mainly composed of silicates (Jessberger *et al.* 2001). Primary candidate materials of interstellar dust are silicate and carbon, although the form of these materials is controversial (Mathis 1990). We cite Mukai (1989) and Hanner (1987) for the complex refractive indices of silicate and carbon, respectively.

Assuming the constituent spherical monomers to be composed of the same material composition in an aggregate, the mass m of the N-monomer aggregate is given by

$$m = \frac{4}{3}\pi s_{\mathrm{m}}^3 N \rho_{\mathrm{m}}, \tag{21}$$

where s_{m} denotes the radius of the identical spherical monomers and ρ_{m} is the bulk mass density of the monomers. The bulk density ρ_{m} of silicate is 2.37×10^3 kg m^{-3}, taken from Lamy (1974), and that of carbon is 1.95×10^3 kg m^{-3}, from the *CRC Handbook of Chemistry and Physics* (Lide 1994).

4.3. Size of Fluffy Dust Aggregates

The accuracy of numerical results from the DDA decreases with increasing size parameter $x_{\mathrm{m}} = 2\pi s_{\mathrm{m}}/\lambda$ of the monomers. While the a1-term method allows us to use a relatively large size parameter of the spherical monomer within the limit $x_{\mathrm{m}} \leq 1$, we safely select the monomer radius $s_{\mathrm{m}} = 0.01$ μm, which yields $x_{\mathrm{m}} \leq 0.45$ (Okamoto 1996, Okamoto and Xu 1998). On the one hand, the average size of monomers building aggregate particles found in interplanetary dust collections is one order of magnitude larger (Brownlee 1985, Jessberger *et al.* 2001). On the other hand, dust particles in the diffuse interstellar medium and dense molecular clouds might be fluffy aggregates consisting of monomers even smaller than $s_{\mathrm{m}} = 0.01$ μm (Mathis and

Whiffen 1989, Ossenkopf 1993). It has been shown that the size of monomers strongly influences the light-scattering properties of aggregates in the visible wavelength range (Levasseur-Regourd *et al*. 1997, Kimura 2001). We therefore become aware that the size of monomers can be an important parameter for estimates of the radiation pressure, to which the radiation pressure cross sections at visible wavelengths largely contribute, but we leave this issue to a future work. The capacity of our computing resources limits the maximum size of the BPCA and BCCA particles to $N = 2048$ monomers. The volume-equivalent radius of the BPCA and BCCA particles defined by $s_{\rm V} = N^{1/3} s_{\rm m}$ ranges over $0.01 \leqq s_{\rm V} \lesssim 0.13$ μm, which corresponds to the mass intervals of $9.9 \times 10^{-21} \lesssim m \lesssim 2.0 \times 10^{-17}$ kg for silicate dust aggregates and $8.2 \times 10^{-21} \lesssim m \lesssim 1.7 \times 10^{-17}$ kg for carbon aggregates.

5. NUMERICAL RESULTS

We introduce $\eta = \sqrt{|\beta_2|^2 + |\beta_3|^2}$ to compare the nonradial component with the radial component of radiation pressure β_1. However, the ratio of nonradial radiation pressure to solar gravity is expressed in terms of $|\beta_3|$ alone, the distribution of which over random orientations of a particle is identical to that of $|\beta_2|$. The numerical values are estimated in 343 orientations for each aggregate with respect to the direction of the incident radiation from the Sun so that randomly oriented aggregates can be simulated.

It should be noted that the cross sections for extinction $C'_{\rm ext}$ and scattering $C'_{\rm sca}$ and the scattering asymmetry parameter $\mathbf{g}'$ are quantities measured in the proper inertial reference frame of the particle. Therefore the β values depend on the orbital velocity $\mathbf{v}$ of the particle through the Lorentz transformation of the wavelengths λ. We demonstrate the velocity effect on the β ratios by providing two examples: the motion of dust particles is either perpendicular or parallel to the direction of propagation of the solar radiation.

5.1. $\mathbf{v} \perp \mathbf{e}_r$

We first ignore the velocity dependence of the β ratios; in other words, we consider the case where a dust particle is in a circular orbit, at perihelion, or at aphelion (i.e., $\mathbf{v} \cdot \mathbf{e}_1 = 0$).

Figures 2 and 3 along the left axis show the ratios of radiation forces to gravitational force on dust grains consisting of silicate and carbon, respectively, as a function of the grain size. The filled triangles and squares are the orientationally averaged values of

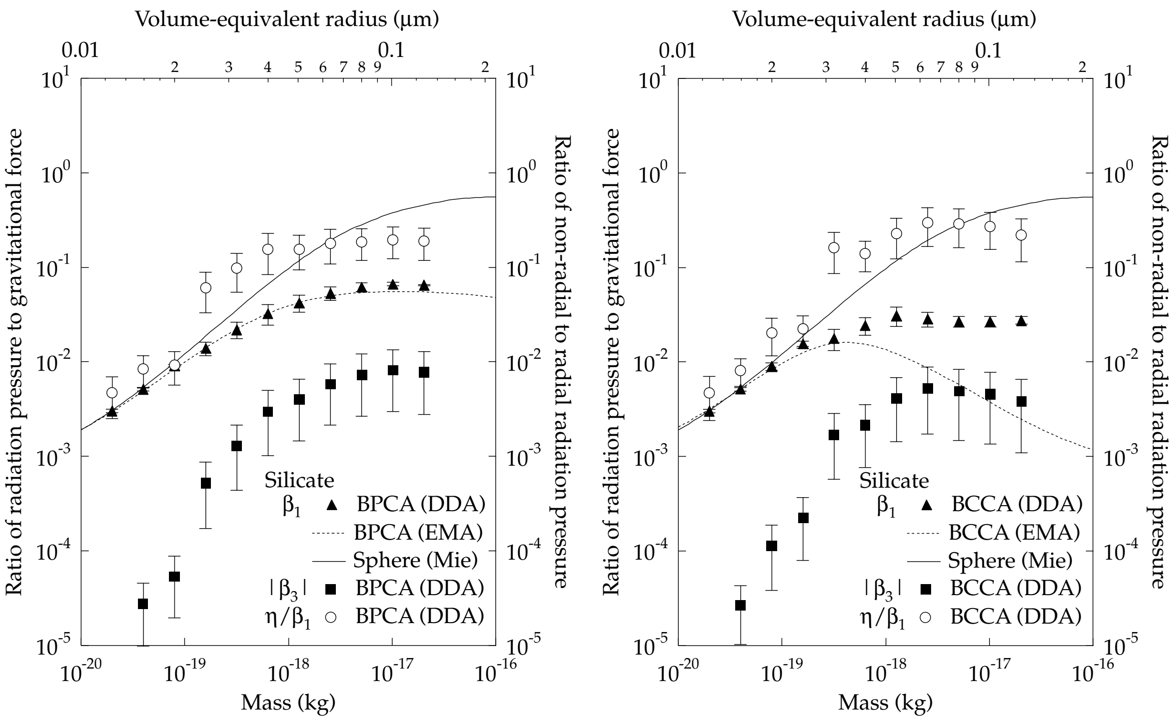

FIG. 2. The ratios of radiation pressure to solar gravitational force acting on the BPCA (left) and BCCA (right) consisting of silicate particles (see the left axis). The top axis indicates the volume-equivalent radius of the aggregates and the bottom axis the mass of the aggregates. The triangles and squares, respectively, are the orientational average values of β_1 and $|\beta_3|$ for aggregate particles calculated by the discrete dipole approximation. The vertical bars denote the standard deviations of the β_1 and $|\beta_3|$, which are evaluated by 343 different orientations of the aggregates with respect to the direction of propagation of incident solar radiation. The solid curves indicate the β_1 values for homogeneous spherical grains evaluated by Mie theory. The dashed lines are rough estimates of the β_1 ratios for aggregate particles using Mie theory combined with the Bruggeman mixing rule. The ratios of nonradial radiation pressure η to radial radiation pressure β_1 are given as open circles (see the right axis).

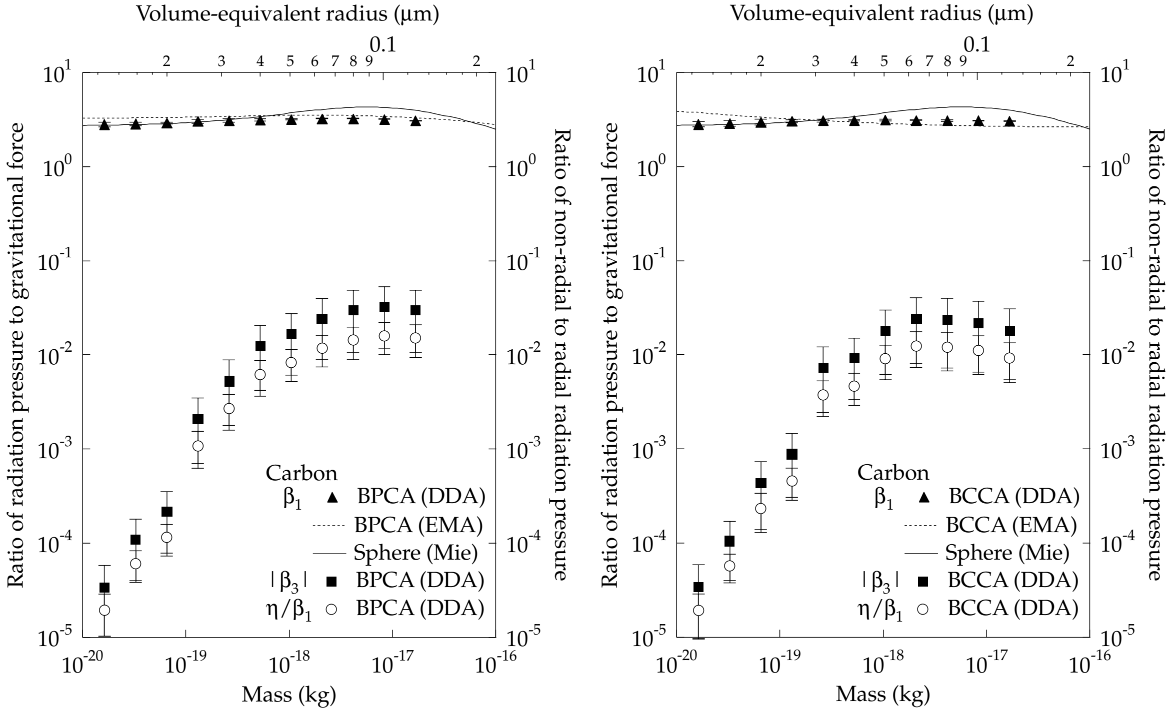

FIG. 3. The same as Fig. 2, but for carbon dust particles.

the β_1 and $|\beta_3|$ for the BPCA (left panel) and BCCA (right panel) particles. The vertical bars indicate the standard deviations of the distributions of β_1 and $|\beta_3|$ originating from the different orientations. The β_1 values of homogeneous spherical grains calculated by Mie theory are included as solid curves, while the $|\beta_3|$ values of the spherical grains are absent because of their symmetrical shape. The dashed lines show rough estimates of the β_1 ratios for the aggregates using the Bruggeman mixing rule combined with Mie theory (Kimura *et al.* 1997). Also inserted in Figs. 2 and 3 along the right axis (open circles) are the ratios η/β_1 of the nonradial to radial radiation pressure.

As the aggregate mass decreases ($N \to 1$), the β_1 ratios for aggregates approach those of homogeneous spherical grains. In comparison with homogeneous spherical grains, the β_1 values of aggregates are less dependent on the aggregate size. As a result, the β_1 values for the aggregates are lower than those of homogeneous spherical grains at submicrometer sizes except for $N = 1$ where the aggregates merge into a spherical constituent particle. Although the EMA used in the calculations fails to reproduce the size dependence of the β_1 ratios for the BCCA particles of silicate, a reasonable agreement can be seen for the BPCA particles.

The $|\beta_3|$ values for the aggregates vanish at $N = 1$, which corresponds to homogeneous spherical grains. This is not the case if the aggregates are composed of nonspherical monomers, although the orientational averaged values become zero. As the number of monomers increases, the $|\beta_3|$ ratios increase but seem to be maximized around a submicrometer radius. This means that the contribution of the nonradial radiation pressure to the dynamics of dust aggregates reaches its peak at submicrometer sizes.

In spite of the difference in shape and structure, our calculations reveal the similarity of the β_1, $|\beta_3|$, and η/β_1 ratios between the BPCAs and the BCCAs, although those values strongly depend on the material composition of the aggregates. The values of β_1 and $|\beta_3|$ show that the radiation pressure acting on carbon aggregates is stronger than that on silicate aggregates in both radial and nonradial directions. However, the η/β_1 values for carbon aggregates are approximately $1/100$ while those of silicate aggregates are on the order of $1/10$. Therefore the nonradial components of the radiation pressure are less important for the dynamics of carbon aggregates compared with silicate aggregates. This results from the fact that absorption of light dominates the radiation pressure cross section for carbon dust particles.

The standard deviation of the β_1 values is not large, implying that the radial radiation force is almost given by the orientationally averaged values irrespective of the particle orientation. In contrast, the large standard deviation of the $|\beta_3|$ values indicates that the magnitude and direction of the nonradial radiation forces largely vary with the orientation of dust aggregates.

5.2. $\mathbf{v} \parallel \mathbf{e}_r$

We next consider the case that the grain velocity is parallel to the direction of propagation of incident solar radiation (i.e., $\mathbf{v} \cdot \mathbf{e}_2 = 0$). This roughly resembles orbits of interstellar dust streaming into the Solar System, β-meteoroids, and dust

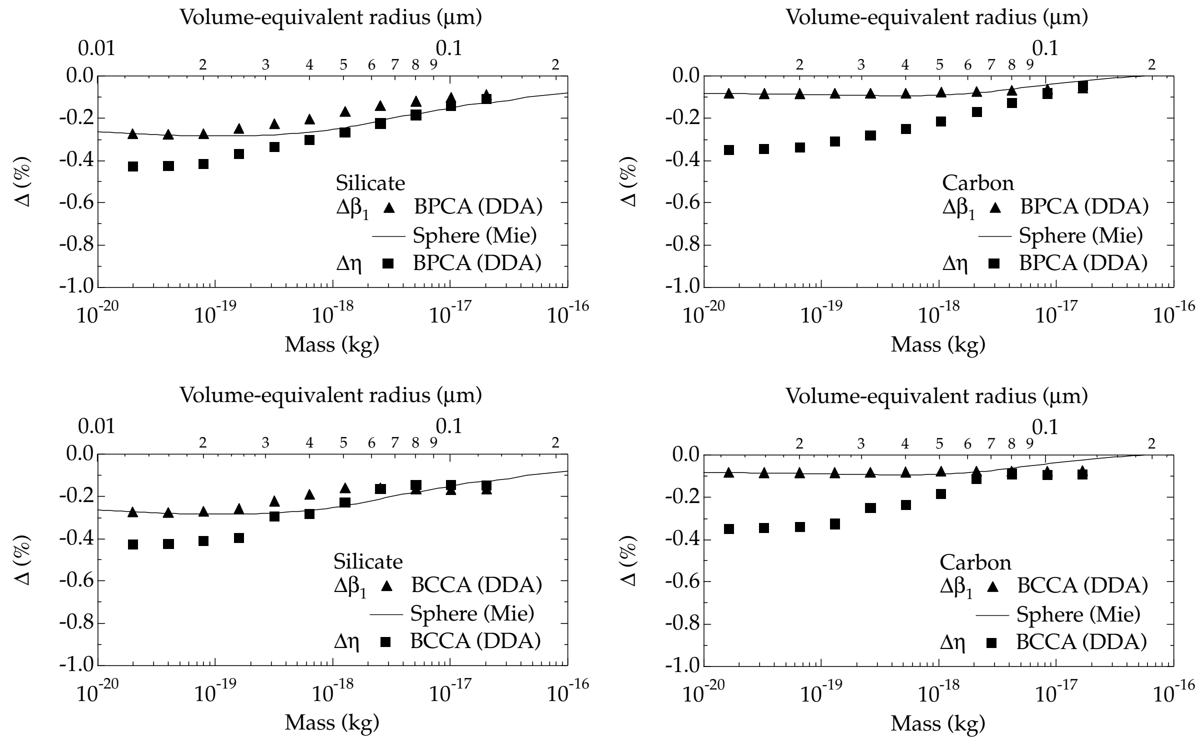

FIG. 4. The deviation Δ of the radiation forces on the BPCA (top) and BCCA (bottom) particles moving radially with $v = 200$ km s^{-1} from those with $v = 0$ km s^{-1}. The triangles and squares are the deviation Δ for the orientationally averaged β_1 and η, respectively, for the aggregate particles calculated by the DDA. The solid curves indicate the deviation Δ of the β_1 for homogeneous spherical grains evaluated by Mie theory. Left and right panels depict the Δ values for silicate and carbon grains in the mass range of 10^{-20}–10^{-16} kg.

particles released from Sun-grazing comets near the perihelion, although in reality they move on complex trajectories (Gustafson and Lederer 1996, Krivov *et al.* 1998, Mann *et al.* 2000). Dust particles released from Sun-grazers may reach $v \sim 200$ km s^{-1} at 1 AU from the Sun, while the heliocentric radial speeds of the β-meteoroids and the interstellar dust grains range from 20 to 30 km s^{-1} (Burns *et al.* 1979, Wehry and Mann 1999, Mann and Kimura 2000). Here we compare the radiation pressure on dust particles moving at $v = 200$ km s^{-1} radially outward from the Sun with that in circular orbits given in Section 5.1. Figure 4 shows the differences in the radiation pressure components $\Delta\beta_1 = (\beta_1^{(v)} - \beta_1^{(0)})/\beta_1^{(0)}$ and $\Delta\eta = (\eta^{(v)} - \eta^{(0)})/\eta^{(0)}$, where the superscripts (v) and (0) indicate that radial speeds of $v = 200$ and 0 km s^{-1}, respectively, are assumed in the computations of the radiation pressure. The filled triangles and squares indicate $\Delta\beta_1$ and $\Delta\eta$, respectively, for the BPCA (top) and BCCA (bottom) particles consisting of silicate (left) and carbon (right). The $\Delta\beta_1$ values for homogeneous spherical grains calculated by Mie theory are inserted as solid curves.

The velocity dependence of radiation pressure is similar for BPCA and BCCA particles and spherical grains. The negative values of $\Delta\beta_1$ and $\Delta\eta$ indicate that the radiation pressure on grains decreases with increasing radial speed of the grains. The dependence of the radial radiation pressure on the grain velocity is weaker than that of the nonradial component. As the radius of dust particles increases from $\sim$0.01 to $\sim$0.1 μm, the nonradial component of radiation pressure on grains moving at high radial speed approaches that in circular orbits. By 1% of accuracy, however, the effect of the grain velocity on the radiation pressure is negligible at least for the parameters considered in this paper.

6. DISCUSSION

6.1. Radial Radiation Pressure

By virtue of the a1-term method of the DDA, we have shown how the β_1 ratio depends on the size, structure, and material composition of dust particles. Numerical results with the DDA a1-term method confirm that the dependence of the β_1 ratios for fluffy aggregates on the size is weaker than that of the β_1 ratios for homogeneous spherical grains. Although a combination of Mie theory and EMAs provides rough estimates for the size dependence of the β_1 ratios except for the silicate BCCAs, a further development of EMAs is desired. In contrast to the porosity dependence of the β_1 ratios estimated by EMAs, the DDA a1-term method reveals that the β_1 ratios are less dependent on the porosity of aggregates while the values strongly vary with the material composition. The solar radiation has therefore a similar dynamical effect for fluffy dust particles of different sizes and porosities if the grains are composed of the same material.

The DDA a1-term method does not allow estimates of radiation pressure acting on aggregates having micrometer or larger radii within our computing resources. The β_1 ratios for large spherical grains are known to be inversely proportional to the dust radius. We therefore expect that the β_1 ratios for fluffy aggregates exceed those of the volume-equivalent spheres in the limit of large particles, because the β_1 ratios for fluffy aggregates show weak size dependences. The β_1 ratios for fluffy dust particles of micrometer or larger sizes could be evaluated using more powerful computers or by other light-scattering theories. One of techniques for computing light-scattering properties of the BPCA and BCCA particles is the superposition T-matrix approach, which is powerful for computing, particularly, orientational averaged properties (Mackowski and Mishchenko 1996, Fuller and Mackowski 1999). An alternative analytical approach for sphere clusters by Xu (1996a,b, 1997, 1998) is superior if huge aggregates are in consideration. This approach can also be used to compute both the radial and nonradial radiation forces on arbitrarily oriented BPCA and BCCA particles. These computing techniques have the advantage of enabling us to study the radiation forces on aggregates whose monomers are larger than the size considered in this paper.

6.2. Nonradial Radiation Pressure

We have shown that the direction of the radiation pressure acting on fluffy dust grains deviates from that of the propagation of the solar radiation. In addition, the force-vector of the Poynting–Robertson effect for fluffy dust particles even in circular orbits is no longer antiparallel to the orbital velocity (see Appendix A). All these differences between fluffy dust particles and homogeneous spherical grains arise from the nonradial component of radiation pressure, which is perpendicular to the direction of propagation of incident radiation. The nonradial radiation pressure on fluffy dust particles of submicrometer sizes can be 1/10–1/100 the radial radiation pressure depending on the material, while the Poynting–Robertson effect is on the order of v/c of the radial radiation pressure (cf. Figs. 2 and 3). Namely, the ratio of the nonradial radiation pressure to the Poynting–Robertson effect can be as large as $0.1c/v$ for weakly absorbing fluffy grains but $0.01c/v$ for strongly absorbing ones. The nonradial component of radiation pressure on submicrometer-size fluffy grains can therefore be greater than the Poynting–Robertson effect, unless their orbital velocities exceed several thousand kilometers per second. If $\beta_2 \cdot v < 0$, then fluffy dust particles can be decelerated by the component of radiation pressure parallel to $\mathbf{e}_2$ rather than by the Poynting–Robertson effect. Otherwise, they can be accelerated by the nonradial radiation pressure when $\beta_2 \cdot v > 0$, instead of being decelerated by the Poynting–Robertson effect.

The orbital inclination of fluffy grains may be altered by the component of radiation pressure parallel to $\mathbf{e}_3$, as discussed by Voshchinnikov and Il'in (1983b), while the inclination of homogeneous spherical grains remains unchanged. On the other hand, the Lorentz force acting on dust grains in the solar magnetic field is known to change the orbital inclination as a consequence of electric charging of dust grains by the solar radiation and solar wind fluxes (Consolmagno 1979, Mukai and Giese 1984, Gustafson and Lederer 1996). The Lorentz forces on homogeneous spherical grains of submicrometer size are estimated at 10 solar radii from the center of the Sun to be 10–100 times smaller than the radial radiation pressure (Mann and Kimura 1997). If this is also the case for fluffy dust particles, the Lorentz force is comparable to the nonradial radiation pressure. The Lorentz force on grains far from the Sun is, however, inversely proportional to the heliocentric distance, in contrast to the inverse square law for solar gravity and radiation pressure (Gustafson 1994). Therefore the Lorentz force is expected to dominate the dynamics of submicrometer-size grains in the outer Solar System. Grün *et al.* (1994) have estimated the Lorentz force on submicrometer-size spherical grains having a surface potential of 5 V at 5 AU from the Sun to be one order of magnitude larger than the solar gravitational force. The Lorentz force on fluffy aggregates might be even stronger if the electric charge on fluffy particles were higher than that on smooth spheres (Mukai 1991). Consequently, the Lorentz force rather than the nonradial components of the radiation pressure may be responsible for changes in the orbital inclination of fluffy grains in the Solar System.

6.3. Rotation of Dust Particles

Whether the nonradial radiation pressure is important for the dynamical evolution of fluffy dust grains depends on the rotation of the grains. The rotation may achieve a certain degree of equilibrium state by bombardment of solar wind particles and solar radiation as well as by the interaction between the electric charges on the grains and the solar magnetic field. The rotation rate and axis of fluffy dust grains may be sensitive to the morphology and composition of the grains, but they are far from being fully understood. Although Draine and Weingartner (1996) have demonstrated that the DDA can be used to estimate the torque on fluffy dust particles induced by solar radiation, such an estimate goes beyond the scope of this paper. When the radiation pressure acting on a fluffy grain is averaged over random orientations at a fixed heliocentric distance, the values of β_2 and β_3 are zero (van de Hulst 1957). However, the ratios β_2 and β_3 need to be averaged over possible orientations according to the rotation of the grain. Furthermore the finite period of the rotation along the trajectory of the grain can prevent the canceling of nonradial radiation forces acting even on a randomly rotating grain. The force-vector of radiation pressure on randomly rotating grains may fluctuate around the direction of incident radiation with time. Similarly the Poynting–Robertson effect on the randomly rotating grains in circular orbits may stagger around the opposite direction of the dust velocity. Therefore fluffy dust particles having a finite period of rotation may take a random walk in the phase space of the orbital elements.

6.4. Lifetime of Dust Particles

Because fluffy grains rotating rapidly in random orientations are weakly perturbed by the nonradial radiation pressure, the equation of motion assuming $\beta_2 = \beta_3 = 0$ can reasonably well describe their trajectories. As a result, the rapidly and randomly rotating grains spiral into the Sun due to the Poynting–Robertson effect even though they are fluffy aggregates. The characteristic time scale of the orbital decay for such a grain is inversely proportional to the β_1 value if the particle is initially in a circular orbit. However, the lifetimes of fluffy dust particles under the Poynting–Robertson effect cannot be approximated by those of spherical grains because of the difference in the β_1 values between fluffy and spherical grains. The lifetimes of fluffy silicate grains having $s_V \approx 0.1\ \mu$m can be one order of magnitude longer than that for spherical grains. The fluffy silicate grains may have longer lifetimes even compared with cylindrical particles because cylindrical grains at $s_V \approx 0.1\ \mu$m have β_1 values larger than spherical grains (Voshchinnikov and Il'in 1983c). Fluffy carbon grains with $s_V \approx 0.1\ \mu$m have lifetimes similar to, but still longer than, those for spherical particles. Large fluffy particles ($s_V > 10\ \mu$m) may be short-lived compared with spherical grains taking into account higher β_1 ratios expected for fluffy grains at the large size limit.

Mukai and Yamamoto (1982) have shown that the Poynting–Robertson effect on homogeneous spherical grains in the submicrometer-size range is comparable to the pseudo-Poynting–Robertson effect, which is caused by solar wind bombardment. It may be straightforward to estimate lifetimes of fluffy dust grains under the pseudo-Poynting–Robertson effect because gas drag forces on fluffy aggregate particles have been computed (Nakamura and Hidaka 1998). It is worthwhile noting that not only the Poynting–Robertson effect and the pseudo-Poynting–Robertson effect but also collisions between dust particles limits their lifetimes (Trulsen and Wilkan 1980, Grün *et al.* 1985, Ishimoto 1998, Ishimoto and Mann 1999). We expect that fluffy dust particles tend to suffer from mutual collisions more often than spherical grains because they have cross-sectional areas that are large compared with spheres of the same volume. In conclusion, particle morphology is one of major factors in determining lifetimes of dust particles in the Solar System under the dynamical and collisional evolutions.

APPENDIX A

Equation of Motion for Arbitrarily Shaped Particles

We derive the equation of motion for fluffy dust particles in the solar radiation field. First two frames of reference are introduced: the solar frame of reference, in which the Sun is at rest and a dust particle has a velocity **v** with respect to the Sun, is denoted by quantities without prime; the particle frame of reference, in which a dust particle is instantaneously at rest, is denoted by primed quantities. We assume that the Sun is a point source of radiation, which is appropriate for dust grains situated at distances large compared with the radius of the Sun (i.e., $r \gg R_\odot$). One may need to take into account the finite size and rotation of the Sun for studying the effect of solar rotation on the dynamics of fluffy dust particles near the Sun as considered previously for spherical dust grains (Guess 1962, Buitrago *et al.* 1989). We further neglect the Yarkovsky effect on the dynamics of dust grains although it may become comparable to the Poynting–Robertson effect for millimeter-to-meter-sized particles (Peterson 1976, Gustafson 1994). A particle is considered to emit isotropically the amount of energy E' absorbed by the particle as thermal radiation in the proper inertial reference frame of the particle; that is,

$$\frac{dE'}{d\tau} = 0, \tag{A1}$$

where the proper time τ is equivalent to the time measured in the proper inertial reference frame. We define the radiation pressure cross section $\hat{C}'_{\mathrm{pr}}$ of a particle by the relation between the momentum $\mathbf{p}'$ delivered to the particle per unit proper time and the incident radiation field,

$$\frac{d\mathbf{p}'}{d\tau} = U'\hat{C}'_{\mathrm{pr}}\mathbf{e}'_r, \tag{A2}$$

where U' is the energy density of the solar radiation at a location of the particle and $\mathbf{e}'_r$ is the direction of propagation of the incident solar radiation. Note that the radiation pressure cross section $\hat{C}'_{\mathrm{pr}}$ is a 3×3 matrix, because $d\mathbf{p}'$ is not necessarily parallel to $\mathbf{e}'_r$.

Equation (A1) and the Lorentz transformations of a four-vector yield (Robertson and Noonan 1968),

$$\frac{dE}{d\tau} = \gamma \mathbf{v} \cdot \frac{d\mathbf{p}'}{d\tau}, \tag{A3}$$

$$\frac{d\mathbf{p}}{d\tau} = \frac{d\mathbf{p}'}{d\tau} + \frac{\gamma^2}{\gamma+1}\left(\frac{\mathbf{v}}{c} \cdot \frac{d\mathbf{p}'}{d\tau}\right)\frac{\mathbf{v}}{c}, \tag{A4}$$

where

$$\gamma = \frac{1}{\sqrt{1-\left(\frac{v}{c}\right)^2}} \tag{A5}$$

is the Lorentz factor and $v = |\mathbf{v}|$ is the speed of the particle with respect to the Sun. The transformation formulas of U' and $\mathbf{e}'_r$ are given by (Robertson and Noonan 1968)

$$U' = w^2 U, \tag{A6}$$

$$\mathbf{e}'_r = \frac{1}{w}\left[\mathbf{e}_r - \frac{\gamma\,(w+1)}{\gamma+1}\frac{\mathbf{v}}{c}\right], \tag{A7}$$

where

$$w = \gamma\left(1 - \frac{\mathbf{v}\cdot\mathbf{e}_r}{c}\right). \tag{A8}$$

If Eqs. (A2), (A6), and (A7) are substituted into Eqs. (A3) and (A4), we obtain the equations of motion in the solar radiation field:

$$\frac{dE}{d\tau} = wUc\left[\gamma\frac{\mathbf{v}\cdot\hat{C}'_{\mathrm{pr}}\mathbf{e}_r}{c} - \frac{\gamma^2\,(w+1)}{\gamma+1}\frac{\mathbf{v}\cdot\hat{C}'_{\mathrm{pr}}\mathbf{v}}{c^2}\right], \tag{A9}$$

$$\begin{aligned}\frac{d\mathbf{p}}{d\tau} = wU\Bigg[&\hat{C}'_{\mathrm{pr}}\mathbf{e}_r - \frac{\gamma\,(w+1)}{\gamma+1}\hat{C}'_{\mathrm{pr}}\frac{\mathbf{v}}{c} + \frac{\gamma^2}{\gamma+1}\frac{\mathbf{v}\cdot\hat{C}'_{\mathrm{pr}}\mathbf{e}_r}{c}\frac{\mathbf{v}}{c}\\ &- \frac{\gamma^3\,(w+1)}{(\gamma+1)^2}\frac{\mathbf{v}\cdot\hat{C}'_{\mathrm{pr}}\mathbf{v}}{c^2}\frac{\mathbf{v}}{c}\Bigg].\end{aligned} \tag{A10}$$

Eq. (A10) coincides with the equation of motion derived by Robertson (1937) when the dust grain is a perfectly absorbing sphere, namely, $\hat{C}'_{\rm pr} = A'\hat{I}$, where A' and $\hat{I}$ denote the geometrical cross section of the spherical particle and the 3×3 unit matrix, respectively.

To terms of order v/c, Eqs. (A9) and (A10) reduce to

$$m\frac{d\mathbf{v}}{dt} = U\left[\left(1 - \frac{\mathbf{v}\cdot\mathbf{e}_r}{c}\right)\hat{C}'_{\rm pr}\mathbf{e}_r - \hat{C}'_{\rm pr}\frac{\mathbf{v}}{c}\right], \tag{A11}$$

where m is the proper mass of the particle. If the dust grain is a perfectly absorbing sphere, Eq. (A11) is identical to the equation of motion derived by Robertson (1937); the first term on the right-hand side exerts a force on a sphere in the direction of incident radiation and the second term exerts a force in the opposite direction of the particle velocity. We call the explicitly velocity-dependent term on the right-hand side the Poynting–Robertson effect and thus the other term is the radiation pressure. In the case of fluffy dust particles, the radiation pressure is parallel to $\hat{C}'_{\rm pr}\mathbf{e}_r$, which is generally not in the direction of incident radiation. Similarly, the Poynting–Robertson effect for fluffy dust particles is also different from that for spherical particles with respect to direction.

Taking into account the solar radiation spectrum, the equation of motion for an arbitrarily shaped particle in the solar radiation field is given by

$$m\frac{d\mathbf{v}}{dt} = \frac{\pi}{c}\left(\frac{R_\odot}{r}\right)^2\int_0^\infty B_\odot\hat{C}'_{\rm pr}\left[\left(1 - \frac{\mathbf{v}\cdot\mathbf{e}_r}{c}\right)\mathbf{e}_r - \frac{\mathbf{v}}{c}\right]d\lambda. \tag{A12}$$

The equation of motion for arbitrarily shaped dust particles has the same form as that for spherical particles except for the matrix description of the radiation pressure cross section. Equation (A12) is equivalent to the expression for a homogeneous spherical particle derived by Soter *et al*. (1977), if $\hat{C}'_{\rm pr} = C'_{\rm pr}\hat{I}$, where $C'_{\rm pr}$ is the usual radiation pressure cross section for a sphere (Bohren and Huffman 1983).

APPENDIX B

Identity of Eqs. (A11) and (10)

We show that Eq. (A11) is identical to Eq. (10) in spite of the difference in the representation of the radiation pressure cross section. To temrs of order v/c, Eqs. (6) and (A2) yield

$$\mathbf{e}'_r\hat{C}'_{\rm pr}\mathbf{e}'_r = C'_{\rm pr}, \tag{B1}$$

$$\mathbf{e}'_\theta\hat{C}'_{\rm pr}\mathbf{e}'_r = -C'_{\rm sca}g'_\theta, \tag{B2}$$

$$\mathbf{e}'_\phi\hat{C}'_{\rm pr}\mathbf{e}'_r = -C'_{\rm sca}g'_\phi. \tag{B3}$$

Inserting Eqs. (1)–(3) into Eqs. (B1)–(B3), we obtain

$$\mathbf{e}_1\hat{C}'_{\rm pr}\mathbf{e}_1 = C'_{\rm pr} - \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}C'_{\rm sca}g'_\theta + \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_1\hat{C}'_{\rm pr}\mathbf{e}_2, \tag{B4}$$

$$\mathbf{e}_2\hat{C}'_{\rm pr}\mathbf{e}_1 = -C'_{\rm sca}g'_\theta + \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_2\hat{C}'_{\rm pr}\mathbf{e}_2 - \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}C'_{\rm pr}, \tag{B5}$$

$$\mathbf{e}_3\hat{C}'_{\rm pr}\mathbf{e}_1 = -C'_{\rm sca}g'_\phi + \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_3\hat{C}'_{\rm pr}\mathbf{e}_2. \tag{B6}$$

Equation (A11) can be written as

$$\begin{aligned} m\frac{d\mathbf{v}}{dt} = U\Bigg\{&\left[\left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right)\mathbf{e}_1\hat{C}'_{\rm pr}\mathbf{e}_1 - \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_1\hat{C}'_{\rm pr}\mathbf{e}_2\right]\mathbf{e}_1 \\ &+ \left[\left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right)\mathbf{e}_2\hat{C}'_{\rm pr}\mathbf{e}_1 - \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_2\hat{C}'_{\rm pr}\mathbf{e}_2\right]\mathbf{e}_2 \\ &+ \left[\left(1 - 2\frac{\mathbf{v}\cdot\mathbf{e}_1}{c}\right)\mathbf{e}_3\hat{C}'_{\rm pr}\mathbf{e}_1 - \frac{\mathbf{v}\cdot\mathbf{e}_2}{c}\mathbf{e}_3\hat{C}'_{\rm pr}\mathbf{e}_2\right]\mathbf{e}_3\Bigg\}, \end{aligned} \tag{B7}$$

where $\mathbf{v}\cdot\mathbf{e}_3 = 0$ is used. If Eqs. (B4)–(B6) are substituted into Eq. (B7), one can finally get Eq. (10).

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