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Creation and expansion of a magnetized plasma bubble for plasma propulsion

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Abstract

We discuss the coupled expansion of a plasma cloud and a neutral gas, both originating from a point-like source located in space and submitted to the action of an external flux of ionizing radiation. This problem is relevant to the artificial magnetospheric propulsion scheme for solar system exploration. We establish the relevant space and time scales for particle diffusion and plasma bubble formation. Emphasis is placed on the low ionization and high collisionallity of the plasma near the point source, leading to the existence of a small ionosphere surrounding the source, where (contrary to the more common views of this propulsion scheme) the ions are not magnetized. The possibility of direct plasma creation by an initial purely neutral gas release is also envisaged.

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1. Introduction

In recent years, a new proposal for spacecraft propulsion was considered, based on the concept of a magnetically confined plasma bubble expansion. This idea, first proposed by Winglee et al. (2000a, b), has been recently renewed by Mendonça et al. (2003) with new concepts. The spacecraft (S/C) would produce a strong magnetic dipole field, and would also contain a plasma source such as a helicon discharge. The resulting plasma bubble would then be pushed by the solar wind. In (Winglee et al., 2000a, b) laboratory testing and computer modelling have demonstrated that the mini-magnetospheric plasma propulsion (M2P2) prototype is able to efficiently produce the plasma, expand the magnetic field lines and create an inflated magnetosphere. In turn, the magnetic force propagation along the flux tubes and acting on the plasma bubble has been evaluated by Mendonça et al. (2003), by averaging over test particle simulations of the dynamic reflection of the solar wind particles on a modified magnetic dipole. Particular solutions for this problem are discussed by Scholer (1970).

The plasma produced by the helicon source should be dense and warm (electron densities and temperatures, n_e and T_e of the order of $5\times10^{13}\,\mathrm{cm}^{-3}$ and 5 eV, respectively). Most probably, the helicon produces a weakly ionized and collisional gas, with $T_e > T_i$, that is injected into the magnetic field generated by the coil. The subsequent evolution depends critically on the values of several physical parameters, some of which do not seem to have been considered in previous work.

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The expected expansion mechanism requires that the charged particles (ions, in particular) are magnetized when they leave the plasma source. However, satisfaction of this condition might not be achieved by simply imposing (via the magnetic field generated by the coil) the width of the plasma source to be a few ion gyroradii. Indeed, in this region, the behavior of the injected gas is also influenced by the collisions of the charged particles that can occur both with the neutral population and among themselves (Coulomb-like). Depending on the relative values of the relevant collision frequencies and the corresponding cyclotron frequencies, several expansion regimes can be anticipated. The extreme cases can be easily pointed out. If the relevant electron and ion collision frequencies are much smaller than the appropriate gyro frequencies then particle magnetization is achieved as discussed by Winglee et al. (2000a, b); and in contrast, occurrence of the inverse situation precludes particle magnetization and the desired expansion mechanism is hindered.

The situation where only one of the charged species becomes magnetized in the source region creates a more physically interesting regime. Assuming that the electrons are the magnetized population and the ions diffuse collisionally, the electrons would move freely along the field lines, but would be transversely tied to them. Because of charge neutrality constraints, both populations are coupled via an electric polarization field; at the outset, the electron gas tends to go ahead of the ions along the field lines and accelerates them through a parallel ambipolar electric field whereas in the perpendicular directions, for sufficiently fast ion diffusion, the generated polarization field would slow down the ions.

The characteristics of the desired expansion of the injected plasma, with the concomitant transport of the coil magnetic field and the ensuing interaction with the solar wind, depend upon the regime that prevails at the source region. It thus seems warranted to study the phenomenology encountered at these early stages of the plasma injection phase.

2. Plasma sources

The inflation of the mini-magnetosphere generated by the injection of plasma into the magnetic field occurs for a non-negligible β-factor of the plasma. The β-factor, defined as the ratio between the plasma pressure $n_e k_B T_e$ and the magnetic field energy $B^2/2\mu_0$, should probably be of the order of 10^{-2} or higher. In the case of a field strength of about 1 kG in the high-field solenoidal region of the dipole field, a β-value of this order implies a plasma with $k_B T_e$ nearly equal to 5 eV and $n_e = 5 \times 10^{13}\,\text{cm}^{-3}$, or a combination thereof with equal product.

The plasma sources able to work in the presence of a strong magnetic field and of producing the high-density plasma required are inductively coupled discharges only. Plasmas generated using electrodes cannot produce such high densities. Among the variety of discharges of this type, helicon discharges seem to be the most promising ones for the purposes of the present application. Helicon discharges are created by an RF-driven antenna in a finite diameter column, in the presence of an axial magnetic field (Chen, 1991; Chen et al., 2001). Helicon sources differ from other inductively coupled plasmas in that a DC magnetic field B is required. The driving frequency is typically 1–50 MHz, with 13.56 MHz commonly used in industrial applications, while the plasma densities range from 10^{11} to $10^{14}\,\text{cm}^{-3}$.

The energy of the RF generator is absorbed by the electrons, due to both collisional and non-collisional mechanisms. It was shown by Chen (1991) that the helicon modes excited in such a device are a superposition of low-frequency whistler modes, with a dispersion relation given by

$$kk_z = \frac{e\mu_0 n_e \omega}{B_z}, \tag{1}$$

with $k = \sqrt{k_z^2 + k_\perp^2}$. In the high-density regime, we have $k_z \gg k_\perp$, so that

$$\lambda_z = \sqrt{\frac{2\pi B_z}{e\mu_0 n_e f}}. \tag{2}$$

In what concerns the energy absorption in the helicon, there is considerable evidence that collisional absorption is too weak to account for energy deposition at low pressures (<10 mTorr) although this mechanism may dominate at higher pressures.

By varying the phase velocity of the wave it is possible to heat electrons by Landau damping, in an electron energy range near the ionization threshold. It follows that it is possible to significantly increase the electron density for the same absorbed RF-power. Experimental evidence of Landau damping has been reported. However, other non-collisional absorption mechanisms may also play a role, such as nonlinear excitation of plasma instabilities or electron acceleration due to the axially non-uniform helicon mode amplitude.

3. Double diffusion process

The existence of a gaseous discharge in space can be seen as a point source of neutral gas and of electron-ion plasma. The expansion of these two components can be studied by using two coupled diffusion processes. Let us first consider the plasma expansion. We consider a magnetic background configuration of the magnetic dipole type, as described in spherical coordinates

(r, θ, ϕ) by

$$\vec{B}(\vec{r}) = \frac{K_0}{r^s}\vec{b}, \quad \vec{b} = 2\cos\phi\vec{e}_r + \sin\phi\vec{e}_\phi, \tag{3}$$

where $K_0 = \mu_0 m/4\pi$, the magnetic moment $\vec{m}$ is oriented along the z-axis, and the field decay exponent s can vary between 1 and 3. For $s = 3$, we have the usual magnetic dipole field. The other values of s are due to the action of the plasma expansion itself on the magnetic dipole field. Also notice that the vector $\vec{b}$ is not a unit vector.

The plasma is described by two fluids, the electron and the ion fluid ($\alpha = \mathrm{e, i}$), according to

$$\frac{\partial n_\alpha}{\partial t} + \nabla \cdot \vec{\Gamma}_\alpha = 0 \tag{4}$$

and

$$\frac{\partial \vec{v}_\alpha}{\partial t} = \frac{q_\alpha}{m_\alpha}(\vec{E} + \vec{v}_\alpha \times \vec{B}) - \frac{S_\alpha^2}{n_\alpha}\nabla n_\alpha - v_\alpha \vec{v}_\alpha, \tag{5}$$

where $S_\alpha = k_B T_\alpha / m_\alpha$ represents the electron and ion thermal velocities. The particle fluxes $\vec{\Gamma}_\alpha$ and the electric field $\vec{E}$ are determined by the usual expressions

$$\vec{\Gamma}_\alpha = n_\alpha \vec{v}_\alpha, \quad \nabla \cdot \vec{E} = \frac{1}{\varepsilon_0}\sum_\alpha q_\alpha n_\alpha. \tag{6}$$

Here, we follow the usual ambipolar approximation, by considering quasi-neutrality $n = n_e \simeq Z n_i$, where Ze is the charge of the ions. We can then derive a diffusion equation in the form

$$\frac{\partial n}{\partial t} - \nabla \cdot \mathbf{D}_a \cdot \nabla n = 0, \tag{7}$$

where the ambipolar diffusion tensor is given by

$$\mathbf{D}_a = \frac{g_e \mathbf{D}_i - g_i \mathbf{D}_e}{g_e - g_i}. \tag{8}$$

In normal ambipolar conditions, we will have $\vec{g} \equiv \vec{E}/E = \nabla n / |n|$. Here, we also used $g_\alpha = \vec{g}\mu_\alpha \vec{E}/E$, where the μ_α is the particle mobility tensor.

If the plasma is diffusing in the presence of an expanding neutral gas, which is being ionized by the incident solar radiation, we have to introduce a source term on the right-hand side of the ambipolar diffusion equation. In some limiting cases (relevant to the plasma bubble formation, as discussed bellow) we can replace the diffusion tensor by a simple radial diffusion coefficient D_a and the ambipolar diffusion equation for the plasma bubble expansion can be written as

$$\frac{\partial n}{\partial t} - \frac{1}{r^2}\frac{\partial}{\partial r}\left(r^2 D_a \frac{\partial n}{\partial r}\right) = v_{ph} n_0(r, t), \tag{9}$$

where v_{ph} is the photoionization frequency, and $n_0(r, t)$ is the density of the expanding neutral gas. In order to completely describe the plasma bubble expansion, we then have to determine the space and time evolution of the neutral gas. Let us assume that n_0 is also described by a diffusion radial process, determined by

$$\frac{\partial n_0}{\partial t} - \frac{1}{r^2}\frac{\partial}{\partial r}\left(r^2 D \frac{\partial n_0}{\partial r}\right) = -v_{ph} n_0 + \delta(r)\delta(t) N_0, \tag{10}$$

where the first source term is associated with the ionization losses and the second represents the injection of N_0 molecules at $r = 0$ and $t = 0$.

In general, the diffusion coefficient for the neutral gas D will vary with the distance r to the source due to the radial variation of the gas density and temperature. Furthermore, it will be coupled to the radial values of the plasma density n due to collisions with the plasma ions. However, as a first approach to the problem we can neglect this collisional coupling and the dependence of D with distance. This is a crude approximation, which will have to be validated later by a more careful analysis. Assuming that such an approximation is valid, we directly used the well-known Gaussian solutions of the type

$$n_0(r, t) = \frac{N_0}{(2\sqrt{\pi D t})^3}\exp(-v_{ph} t - r^2/4Dt). \tag{11}$$

This is valid for a release of N_0 neutral atoms at time $t = 0$. This result can be generalized to the case where the total number of neutral atoms N_0 is released at a rate R_0 (number of atoms per second) during a time interval Δt. The corresponding solution for the density of neutrals is now determined by

$$n_0(r, t) = \int_0^{\Delta t} \frac{R_0}{(2\sqrt{\pi D (t - t')})^3} \times \exp\left[-v_{ph} t - \frac{r^2}{4D(t - t')}\right] dt', \tag{12}$$

where, obviously the rate of neutral emission by the source located at $r = 0$ is $R_0 = N_0/\Delta t$. We can now use this solution in the ambipolar diffusion equation in order to establish the law for the plasma bubble formation. The result is

$$n(r, t) = \int_0^t dt' \int_0^\infty dr' \, S(r', t') G(r, r' | t, t') + n_{free}(r, t), \tag{13}$$

where the source term for the plasma density is due to the ionization of neutrals $S(r, t) = v_{ph} n_0(r, t)$, and n_{free} is the solution resulting from the free ambipolar diffusion of the initial plasma density $n(r, t = 0)$. For a release of plasma by the spacecraft at $t = 0$, this free solution would be formally identical to Eq. (11), with D replaced by the ambipolar diffusion D_a and N_0 by the initial number of released plasma particles $N_p(0)$. In order to determine the forced term in Eq. (13) we have to make use of the well-known Green's function for the diffusion equation

$$G(r,r'|t,t') = \frac{H(t-t')}{(2\sqrt{\pi D_a(t-t')})^3} \exp\left[-\frac{(r-r')^2}{4D_a(t-t')}\right], \quad (14)$$

where $H(t-t')$ is the Heaviside or step function. Solution (13) will allow us to compare the ionization contributions for the plasma bubble formation, with the free expansion of the plasma produced by the plasma discharge. Noting that the number of neutrals launched by the discharge is nearly two orders of magnitude larger than the number of ionized particles, $N_0 \gg N_p(0)$, we can conclude that, for sufficiently high values of the ionization frequency v_{ph}, the first term will eventually become dominant.

This formulation gives a general picture of how the coupled expansion of plasma and molecular gas has to be treated. However, its explicit solution constitutes a difficult task that can only be achieved by numerical integration of the coupled diffusion equations, since both the diffusion coefficient of neutral gas and the ambipolar diffusion coefficient depend on the gas density and on the electron density, respectively. This direction of research will be pursued in the future.

4. Photoionization of the neutral gas

The available helicon plasma sources produce a weakly ionized plasma, with a degree of ionization less than 1%. This means that the plasma source is actually launching a nearly neutral gas into space. It is then interesting to study the photoionization of the injected neutrals by the solar wind and the solar UV radiation, in particular. This study will also allow us to raise the possibility of creating the mini-magnetosphere used for plasma propulsion without the need of a plasma source. If possible, this plasma bubble creation by photoionization would allow for a considerable simplification of the artificial magnetospheric propulsion scheme. Notice that, as soon as the gas injected into space gets ionized, it will become trapped by the magnetic dipole created by the spacecraft coils. In that sense this configuration is clearly distinct from the ionospheric release experiments done in the past, because of the local magnetic confinement.

In order to evaluate plasma production by direct photoionization, we recall that the average irradiance of the sun E_0, assumed as a blackbody source, is proportional to the fourth power of the temperature $E_0 = \sigma_S T^4$, where σ_S is the well-known Stefan–Boltzmann constant. Thus, for an average temperature of the surface of the sun given by $T = 5780\,\mathrm{K}$ we obtain the irradiance value of $E_0 = 6.33 \times 10^7\,\mathrm{W\,m^{-2}}$.

The resulting energy received on Earth, or at a mean distance of $R = 1.5 \times 10^8\,\mathrm{km}$ from the sun, will be $E = E_0 \times (r/R)^2 = 1.38 \times 10^3\,\mathrm{W\,m^{-2}}$, where $r = 7 \times 10^5\,\mathrm{km}$ is the radius of the sun. The efficiency of this nearly blackbody radiation to photo-ionize a given neutral gas depends very much on the position of its peak value with respect to the ionization threshold of the gas. According to the well-known Wien displacement law, the wavelength of peak emission for the sunlight is given by $\lambda_{max}(\mu\mathrm{m}) = 2897/T$ (K). Thus, for $T = 5780\,\mathrm{K}$, we obtain $\lambda_{max} = 5010\,\mathring{A}$. For instance, the ionization threshold in argon

$$hv + \mathrm{Ar} \longrightarrow \mathrm{Ar}^+ + \mathrm{e} \quad (15)$$

is equal to $E_{ion} = 15.76\,\mathrm{eV}$, which corresponds to an ionization wavelength of $\lambda^* = hc/E_{ion} = 787\,\mathring{A}$, deep in the ultraviolet. So, the ratio of photon energy to thermal energy for this ionization threshold is $\eta = hc/(\lambda^* k_B T) = 31.63$. This means that we can take the ultraviolet approximation $hc/\lambda \gg k_B T$ of the Planck distribution and write, for the mean spectral energy density of photons

$$u_\lambda \simeq \frac{8\pi hc}{\lambda^5} \exp\left(-\frac{hc}{\lambda k_B T}\right). \quad (16)$$

On the other hand, using the mean energy density

$$u_0 = \int_0^\infty u_\lambda \, \mathrm{d}\lambda = \frac{8\pi^5}{15} \frac{(k_B T)^4}{(hc)^3}, \quad (17)$$

we obtain, for photons of wavelength $hc/\lambda \gg k_B T$, an energy distribution given by

$$u_\lambda \simeq u_0 \frac{15}{\pi^4} \frac{\eta^4}{\lambda} \mathrm{e}^{-\eta}. \quad (18)$$

Assuming now that the most significant part of the integral over the energy spectrum corresponds to energies above the ionization threshold, $E \geq E_{ion} = 15.76\,\mathrm{eV}$ (i.e. for $\lambda \leq \lambda^* = 787\,\mathring{A}$), we obtain a fractional ration for ionizing photons given by

$$F = \frac{1}{u_0} \int_0^{\lambda^*} u_{\lambda^*} \, \mathrm{d}\lambda^* \simeq \frac{u_\lambda \lambda^*}{u_0} = \frac{15}{\pi^4} \eta^4 \mathrm{e}^{-\eta} = 2.8 \times 10^{-9}. \quad (19)$$

The irradiance from the sun, in the short wavelength range $\lambda \leq \lambda^* = 787\,\mathring{A}$, will then be given by

$$I(R) = \frac{E}{R^2} F = \frac{3.86 \times 10^{-6}}{R^2}\,\mathrm{W\,m^{-2}} \quad (20)$$

with R in AU (distance from the center of the sun to the center of the earth). We can also determine the ionizing photon flux as

$$\Gamma(R) = \frac{I(R)}{hv} = \frac{1.53}{R^2} \times 10^{12}\ \mathrm{photons\,s^{-1}\,m^{-2}} \quad (21)$$

to which corresponds a photoionization frequency of

$$\nu_{ph}(R) = \Gamma(R)\sigma = \frac{3.22}{R^2} \times 10^{-9}\,\mathrm{s}^{-1}, \tag{22}$$

where $\sigma \simeq 21.06$, with $1\,\mathrm{Mb} = 10^{-18}\,\mathrm{cm}^2$, is the photoionization cross section for Argon (Verner et al., 1993; Verner and Yakovlev, 1995). We know that, for Argon, at a pressure $p = 760\,\mathrm{Torr}$, temperature $T_0 = 273\,\mathrm{K}$, and density $n_L = 2.69 \times 10^{19}\,\mathrm{cm}^{-3}$, we have a diffusion coefficient for the neutral atoms of $D^* = 0.157\,\mathrm{cm}^2\,\mathrm{s}^{-1}$. So, if the neutral Argon is released by the S/C at a pressure of $p = 1\,\mathrm{Torr}$ and temperature $T_0 = 300\,\mathrm{K}$, we will have a diffusion coefficient of $D = n_L D^*/n_0 = 131\,\mathrm{cm}^2\,\mathrm{s}^{-1}$, where $n_0 = 3.22 \times 10^{16}\,\mathrm{cm}^{-3}$ is the number of molecules injected around $r = 0$ per unit volume.

These values can be used to determine the neutral gas cloud formed around the spacecraft, as described in the previous section. The average distance travelled by one neutral atom in a time t will be determined by $l = \sqrt{3\langle x^2 \rangle}$, where the mean-square displacement $\langle x^2 \rangle$ is

$$\langle x^2 \rangle = \frac{1}{N_0} \int x^2 n(\vec{r}, t)\, \mathrm{d}\vec{r} = 2Dt, \tag{23}$$

so that $l = \sqrt{6\,Dt} = \sqrt{6 \times 131\, t}$, or $l(\mathrm{cm}) = 28 \times \sqrt{t(\mathrm{s})}$.

While expanding, this neutral gas becomes photoionized by the incident sunlight.

5. Electron diffusion losses

Once the neutral atoms are ionized, their transport properties change considerably, because the resulting charged particles become trapped by the dipole magnetic field. Diffusion will become ambipolar, with the electrons moving faster and being slowed down by the ions. Let us consider the ambipolar diffusion of the electrons in the nearly infinite and nearly empty medium which surrounds the spacecraft (S/C). Such a diffusion takes place in the presence of the ambipolar space-charge field, and of the magnetic field created by the S/C. Let us consider diffusion directed along some radial direction with respect to the S/C. If diffusion is taking place parallel to the magnetic field, $\vec{B}$, we can write

$$D_{a\parallel} = \frac{D_e \mu_i + D_i \mu_e}{\mu_e + \mu_i} \simeq \mu_i \frac{KT_e}{e}, \tag{24}$$

where

$$D_\alpha = D_{\alpha\parallel} = \frac{KT_\alpha}{\mu^*_{\alpha 0} \nu_{\alpha 0}}, \quad \mu_\alpha = \mu_{\alpha\parallel} = \frac{e}{\mu^*_{\alpha 0} \nu_{\alpha 0}} \tag{25}$$

with $\alpha = (\mathrm{e}, \mathrm{i})$ for electrons and ions. Here, $\mu^*_{\alpha 0}$ is the reduced mass for collisions with neutrals ($\mu^*_{e0} = m_e$ and $\mu^*_{i0} = M_i/2$).

In Argon, we can use $\mu_{i0} = 1.9\,\mathrm{cm}^2\,\mathrm{V}^{-1}\,\mathrm{s}^{-1}$, and $\mu_i = \mu_{i0} \times (n_L/n_0) = 0.159\,\mathrm{m}^2\,\mathrm{V}^{-1}\mathrm{s}^{-1}$, for the above values of n_L and n_0. If we further assume $k_B T_e/e = 5\,\mathrm{V}$, we obtain $D_{a\parallel} = 0.795\,\mathrm{m}^2\,\mathrm{s}^{-1}$.

If, on the contrary, electron ambipolar diffusion is taking place in a direction perpendicular to the magnetic field, the diffusion coefficient will be determined by

$$D_{a\perp} = \frac{D_{e\perp}\mu_{i\perp} + D_{i\perp}\mu_{e\perp}}{\mu_{e\perp} + \mu_{i\perp}} \tag{26}$$

with

$$D_{\alpha\perp} = D_{\alpha\parallel} \frac{1}{1 + \mu^2_{\alpha\parallel} B^2}, \quad \mu_{\alpha\perp} = \mu_{\alpha\parallel} \frac{1}{1 + \mu^2_{\alpha\parallel} B^2}, \tag{27}$$

for $\alpha = (\mathrm{e}, \mathrm{i})$, and $D_{\alpha\parallel}$ and $\mu_{\alpha\parallel}$ given in Eq. (25). Now inserting Eq. (27) into the expression for $D_{a\perp}$, we obtain (Golant et al., 1980)

$$D_{a\perp} = D_{a\parallel} \frac{1}{1 + \mu_{i\parallel}\mu_{e\parallel} B^2}. \tag{28}$$

For a magnetic field amplitude of $B = 1\,\mathrm{kG}$ (as expected near the S/C), and taking into account that $\mu_e \simeq 235 \times \mu_i$ in Argon, we can derive the following relation between the parallel and the perpendicular diffusion coefficients:

$$D_{a\perp} = D_{a\parallel} \frac{1}{1 + 5.94 \times 10^{-2}} = 0.94\, D_{a\parallel}. \tag{29}$$

This means that, even if the electrons are strongly magnetized (in the sense that their collision frequencies are much lower than the electron cyclotron frequency), no significant effect associated with the magnetic field occurs in the diffusion process, which will then remain rather spherical around the S/C, as considered before. This is true, at least for the early stages of the diffusion process, when the neutral gas density is still quite high, of the order of $n_0 = 3.22 \times 10^{16}\,\mathrm{cm}^{-3}$. For later stages, when the neutral gas pressure goes down to $p = 0.1\,\mathrm{Torr}$, we will have $n_0 = 3.22 \times 10^{15}\,\mathrm{cm}^{-3}$ and magnetization effects will start to dominate over collisions with the neutrals. In this case, we will obtain $D_{a\perp} = 0.144 D_{a\parallel}$. The electron cloud (and by ambipolarity, the plasma cloud itself) will then start to deviate from spherical diffusion. However, at the later stages of expansion, the electron density goes down to such low levels that the ambipolar diffusion model is no longer appropriate, and the smooth transition from ambipolar to free electron diffusion has to be considered. Once again, only numerical solutions can eventually be given, since the electron diffusion is strongly coupled to the expansion of neutral gas. This question is the focus of the present work.

6. Photoionization processes

We have considered above the specific example of Argon. Other gases can eventually be used to create the plasma bubble, with more favorable ionization energies and photoionization cross-sections. A few cases are given in the following table. These few examples show that it is possible to envisage the plasma bubble creation without using any plasma discharge and just relying on the photoionization processes of neutral gas release.

	Z	E_{ion} (eV)	σ (Mb)	ν_{ph} (s^{-1})
H	1	13.60	5.475×10^4	4.11×10^{-4}
He	2	24.59	9.492×10^2	1.10×10^{-14}
Li	3	5.39	6.245×10	4.21×10^{-1}
N	7	14.53	8.235×10^2	1.17×10^{-6}
O	8	13.62	1.745×10^3	1.26×10^{-5}
Ar	18	15.76	2.106×10	3.22×10^{-9}

7. Conclusions

We have examined the processes of plasma bubble formation and expansion around a magnetized spacecraft (S/C), as proposed for the artificial magnetospheric propulsion scheme. These processes can be described by two coupled equations: one for the neutrals and the other for the ambipolar plasma expansion. Typical length and time-scales were established.

Creation of a plasma bubble by photoionization of the neutral gas released by the S/C is, in principle, an interesting alternative to be considered for this propulsion scheme. This is due to the low degree of ionization of the plasma created by typical plasma sources, such as the helicon sources, installed inside the S/C. If possible, this would considerably reduce the weight and the energy needs, and simplify the propulsion scheme. As shown here, photoionization of Argon gas by solar UV radiation does not seem to be very efficient because of the high ionization threshold. However, higher photoionization rates can be obtained using other gases (i.e. Hydrogen or Lithium), as shown in the above table.

The ionization by impact of electrons and ions from the solar wind ($n_e = n_i = 5 \times 10^6\,\text{m}^{-3}$, $v_e = v_i = 400\,\text{km}\,\text{s}^{-1}$) could eventually contribute to the plasma bubble formation, but a simple estimate shows that their influence is negligible. Of course, if this simple scheme of plasma formation fails, there is always the possibility of using a plasma source, such as the helicon discharge, according to the original proposal for the M2P2. Future studies will require a more elaborate description of the plasma bubble formation, with non-spherical solutions for the ambipolar expansion.

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