4. V. I. Davydenko, G. I. Dimov, and G. V. Roslyakov, "Obtaining precision high-intensity ion and atomic beams," Dokl. Akad. Nauk SSSR, 271, No. 6 (1983).
5. V. I. Batkin, V. N. Getmanov, I. M. Ikryanov, and O. Ya. Savchenko, "Proton source with a regulated current for an electrostatic accelerator," Vopr. At. Nauki Tekh. Ser. Tekh. Fiz. Eksp., No. 1/22, Kharkov (1985).
6. V. N. Getmanov and O. Ya. Savchenko, "Pulsed proton linear accelerator," Otkryt. Izobret., No. 26 (1989).
7. V. I. Davydenko, I. I. Morozov, and G. V. Roslyakov, "Proton source of the hydrogen-atom injector of the AMBAL device," Preprint No. 85-3 [in Russian], Akad. Nauk SSSR, Sib. Otd., Inst. Yad. Fiz., Novosibirsk (1980).
8. V. I. Batkin and O. Ya. Savchenko, "Effect of a weak magnetic field in the expander of an ion source for controlling the beam current and profile," Prib. Tekh. Éksp., No. 3 (1989).
9. V. N. Getmanov and O. Ya. Savchenko, "Effect of a weak magnetic field on the value of the current of an arc proton source," Zh. Prikl. Mekh. Tekh. Fiz., No. 3 (1987).
10. V. I. Batkin, "Beam modulation by a grid electrode in the expander of an ion source," in: Author's Abstracts, 11-th All Union Seminar on Linear Accelerators of Charged Particles [in Russian], Kharkov (1989).

FARADAY DARK SPACE FOR A DISCHARGE IN HELIUM

A. I. Ivanchenko and A. É. Medvedev UDC 537.525

Discharges in monatomic gases have been researched for more than half a century, but our knowledge of the Faraday dark space is still inadequate [1]. Existing models give a satisfactory description only at pressures up to a few torr [2, 3]. In experiments [4] with helium at 1.3-9.3 kPa (10-70 torr), the length of the dark space was many times the tube radius. The result [4] is unexpected from the classical viewpoint [3] but agrees qualitatively and quantitatively with the model proposed here.

We consider the transition region from the weak-field area in the cathode layer to the positive column. If the plasma is quasineutral, the electron temperature is constant over the cross section of the tube, and the gas heating is slight [1, 3], so we get the solution for that region from the particle conservation equation

$$\nabla\left(D\nabla n-\frac{\mathbf{j}}{e}\frac{\mu_+}{\mu_e(E)}\right)=\nu(E)\,n-\beta n^2, \tag{1}$$

in which D, μ_+, μ_e, β, ν are the ambipolar diffusion coefficient, the ionic and electron mobilities, the bulk recombination constant, and the ionization frequency. On the left in (1), the ambipolar drift accompanies the diffusion one, which is due to the electron mobility being dependent on the field [5, 6]. We neglect the generation and loss terms for the bulk in the Faraday dark space. The particle balance in the positive column is governed not only by plasma diffusion to the wall but also by bulk processes, and therefore the proportion of ionization falls sharply as one passes from the positive column to the Faraday dark space because the ionization term is exponentially dependent on the field, and it cannot balance the plasma loss to the wall. For example, if the field in helium is reduced by 10% from that in the positive column, the ionization proportion becomes less than a third of the diffusion loss, and as the field falls further, it ceases to play a part in the balance. A detailed comparison can be made for the recombination losses and the diffusion ones from the [7, 8] results and the [9] temperature dependence; incorporating the bulk recombination term is necessary for electron temperatures less than 0.1-0.16 eV (for p = 1.3-9.3 kPa and j = 10^{-2} A/cm^2). The rates of the processes are similar in magnitude in each branch of the collisional-radiative recombination involving electrons. The dissociative recombination rate

Novosibirsk. Translated from Zhurnal Prikladnoi Mekhaniki i Tekhnicheskoi Fiziki, No. 1, pp. 12-14, January-February, 1991. Original article submitted July 14, 1989.

at 1.33 kPa is relatively low. At positive-column temperatures, the bulk recombination losses are larger than the diffusion ones for the various inert gases at $p \geq 6.5$-13 kPa.

Then we write (1) without the loss and generation terms for the bulk and with transport coefficients independent of pressure:

$$\nabla\left(\frac{Dp}{p}\nabla n - \frac{\mathbf{j}}{e}\frac{\mu_+ p}{\mu_e p}\right) = 0. \tag{2}$$

We consider (2) for a tube with x as the axis for two limiting cases: low and high pressures. At low pressures, we neglect the last term and get the standard solution with characteristic dimension along x of the order of the tube radius [3], in which the plasma loss from the cathode region by ambipolar diffusion is balanced by the diffusion transport to the walls. As the pressure rises, the drift loss from the cathode layer begins to predominate over the diffusion loss because the electron mobility is dependent on the field. We estimate the pressure at which the diffusion term becomes unimportant.

The diffusion component of the current is small up to $E/p = 0.03$-0.05 V/cm·torr, so the electron density is $n \approx j/e\mu_e E$. In helium [1, 10]

$$D \approx \mu_+ T \approx k\mu_+ E/p \quad (k \approx 2 \text{ torr·cm}), \tag{3}$$

and the ambipolar diffusion term is

$$F = \frac{d}{dx} D \frac{d}{dx} n = 2\mu_+ \frac{j}{e} \frac{d}{dx} \frac{(-E/p)}{(\mu_e E)^2} \frac{d}{dx}(\mu_e E). \tag{4}$$

We approximate the field dependence of the electron mobility at 0.1-1 V/cm·torr from [11] as

$$\mu_e p = A + \frac{B}{E/p}, \quad A \approx 8.5\cdot 10^5 \text{ cm·torr/(V·sec)}, \quad B \approx 1.6\cdot 10^5 \text{ cm/sec} \tag{5}$$

and substitute into (4) to get

$$F = \left[\frac{(E/p - B/A)}{(E/p + B/A)}\left(\frac{dE/p}{dx}\right)^2 - \frac{E}{p}\frac{d^2}{dx^2}\frac{E}{p}\right]^2 \frac{\mu_+ j/e}{A(E/p + B/A)^2}. \tag{6}$$

From (5), the ambipolar drift term is

$$G = \frac{j}{e}\mu_+ p \frac{d}{dx}\frac{1}{\mu_e p} = \frac{j}{e}\mu_+ p \frac{B}{(AE/p + B)^2}\frac{dE/p}{dx}. \tag{7}$$

Estimates are $d(E/p)/dx \approx (E/p)_{pc}/\Delta x$, $d^2(E/p)/dx^2 \approx (E/p)_{pc}/(\Delta x)^2$, in which $(E/p)_{pc}$ is E/p in the positive column, while Δx is the characteristic scale of the inhomogeneity along the tube axis, so (6) and (7) give

$$\max\left\{\frac{F}{G}\right\} = \frac{\left[\frac{(E/p)_{pc}}{\Delta x}\right]^2}{\frac{p}{2}\frac{B}{A}\frac{(E/p)_{pc}}{\Delta x}} \approx \frac{(E/p)_{pc}}{0.1 p\Delta x}. \tag{8}$$

The diffusion-limited solution ($\Delta x \approx R$, R tube radius) ceases to apply if $pR \gtrsim 10$ torr·cm. For ambipolar drift, $\Delta x \approx 0.1\, pR^2$, and the ambipolar diffusion term falls even more rapidly as the pressure rises: $\max\{F/G\} \approx (E/p)_{pc}/(0.1pR)^2$.

These estimates can be used with the balance equation, the boundary conditions, and the expression for the current density, which consists only of the drift component for the electron current j_x along the axis, on the assumption of a quasihomogeneous solution such that $dj_x/dx \ll j/L$ (L is the scale of the transition region) to get a treatment with separable variables:

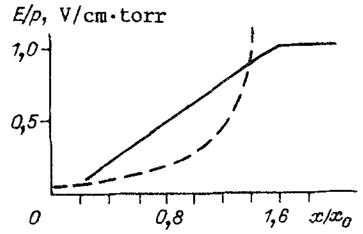

Fig. 1

$$\frac{1}{r}\frac{d}{dr}D(E/p)\,r\frac{dn}{dr}+\frac{j_x}{e}\mu_+\frac{d}{dx}\frac{1}{\mu_e(E/p)}=0,$$
$$|\mathbf{j}|\approx j_x=en\mu_eE,\quad 2\pi\int_0^R j_x r\,dr=I, \tag{9}$$
$$E/p|_{x=0}=(E/p)_0,\ n|_{r=R}=0$$

(I is the total discharge current). We divide the balance equation by D and substitute n = Z(x) × Y(r) into (9) to get

$$Y(r)=J\left(r\sqrt{\frac{E}{T}\nabla_x\left[\ln\frac{1}{\mu_e}\right]}\right). \tag{10}$$

As in Schottky's theory [12], we assume that the boundary condition $n|_{r=R}=0$ is equivalent to the condition for the first zero in the zero-order Bessel function (Y(R) = 0); from (10), this gives a solution along the axis

$$dx=\left(\frac{R}{2.4}\right)^2 p\,\frac{E/p}{T}\,d\left[\ln\frac{1}{\mu_e p}\right]. \tag{11}$$

We integrate (11) with upper limit defined by emergence from the positive column and use (3) to get

$$\int_0^L dx=\left(\frac{R}{2.4}\right)^2\frac{p}{2}\ln\frac{1}{p\mu_e(E/p)}\Big|_{(E/p)_0}^{(E/p)_{\text{pc}}}. \tag{12}$$

The lower integration limit (x = 0 in the cathode region) is determined by the applicability bound for this model. As $(E/p)_0$ we took 0.03 V/cm·torr, below which one cannot neglect recombination in the bulk and the diffusion component in the electron current. Finally, as we know that for helium with pR ≳ 10 torr·cm we have $(E/p)_{\text{pc}}\approx 1$ V/cm·torr [10], we have the scale of the transition region as

$$L=pR^2\cdot 0.087\ln\frac{4\cdot 10^6}{10^6}\approx 0.12pR^2. \tag{13}$$

We use the analytic form $\mu_e p=\mu_e p(E/p)$ in the range E/p = 0.01-1 V/cm·torr with the [11] data in the form

$$\mu_e p=A+\frac{B}{E/p+C}, \tag{14}$$

in which $A = 8\cdot 10^5$ cm²·torr/V·sec, $B = 1.9\cdot 10^5$ cm/sec, and $C = 2.5\cdot 10^{-2}$ V/cm·torr to get the field distribution in the Faraday dark space. We substitute (14) into (12) and invert the dependence for E/p to get (Fig. 1, dashed line)

$$E/p(x)=\left[\left(\frac{A}{B}+\frac{1}{(E/p)_0+C}\right)\exp\left(-\frac{x}{x_0}\right)-\frac{A}{B}\right]^{-1}-C,\quad x_0\approx 0.087pR^2.$$

Experiment [4] in our pressure range is represented by the solid line in Fig. 1, with the scale of the Faraday dark space given by $pL = a(pR)^2$, $a \approx 0.14\ (\text{torr}\cdot\text{cm})^{-1}$. The proportionality observed between L and pR^2 [compare (13)] is governed by the dependence of F on p and R, with G independent of those parameters [see (9)]. It is not surprising that there is an anomalously large dark space only in helium [4] because already for Ne and particularly for the heavier gases, the lower pressure bound for the model increases in proportion to T/(E/p), as (11) shows, from 10 torr for helium (R = 2 cm, nominal boundary for model L ≈ 2.5R) to more than 50 torr, where bulk recombination becomes considerable for the inert gases beginning with Ne.

LITERATURE CITED

1. Yu. P. Raizer, Gas-Discharge Physics [in Russian], Nauka, Moscow (1987).
2. Yu. M. Kagan, C. Cogen, and P. Avivi, "Faraday dark space of a He glow discharge," J. Appl. Phys., 63, No. 1 (1988).
3. V. A. Granovskii, Electric Currents in Gases: Steady-State Current [in Russian], Nauka, Moscow (1971).
4. A. Rutcher and Ch. Wilke, "Investigation of very long Faraday dark spaces," 11th Intern. Conf. on Phenomena in Ionized Gases, Prague (1973), contr. pap.
5. D. I. Vysikailo, "Ambipolar drift," Teplofiz. Vys. Temp., 23, No. 4 (1985).
6. Yu. S. Akishev, D. I. Vysikailo, A. P. Napartovich, and V. V. Ponomarenko, "A study on a quasistationary discharge in nitrogen," Teplofiz. Vys. Temp., 18, No. 2 (1980).
7. J. Berlande, M. Cheret, R. Deloche, et al., "Pressure and electron density dependence of the electron—ion recombination coefficient in helium," Phys. Rev., 1A, No. 3 (1970).
8. R. Deloche, A. Confalone, and M. Cheret, "Mesure en fonction de la pression, de la variation du coefficient de recombination electron-ion, dans l'helium faiblement ionisé, en regime de relaxation," C. R. Acad. Sci., 267, No. 18 (1968).
9. B. M. Smirnov, Atomic Collisions and Elementary Processes in Plasmas [in Russian], Nauka, Moscow (1988).
10. A. Van Engel, Ionized Gases [Russian translation], Fizmatgiz, Moscow (1959).
11. I. McDaniel, Collision Processes in Ionized Gases [Russian translation], Mir, Moscow (1967).
12. W. Schottky, "Diffusions-Theorie der positiven Säul," Phys. Z., 25, No. 23 (1924).

MHD DECELERATION AND HEAT TRANSFER FOR A SPHERE IN A SUPERSONIC FLOW OF PARTIALLY IONIZED GAS

V. V. Gubin and V. A. Shuvalov

UDC 533.95:538.4:537.523

The inherent magnetic field has a substantial effect on the structure of the perturbed zone at the surface of a body in a supersonic low-density plasma [1]. The inherent field of the body may be due to a set of currents or to permanent magnets. The perturbations caused by that field affect the functional and dynamic characteristics in the interaction with the flow. An approximate numerical analysis [2, 3] and experiment [4, 5] indicate effective MHD retardation in such a flow. It is desirable to examine MHD control of heat transfer and the aerodynamic characteristics.

Measurements are reported here on the MHD retardation and heat transfer as affected by the direction of the inherent field of the body $\mathbf{H}$ with respect to the incident velocity vector $\mathbf{U}_\infty$. MHD control is possible for the aerodynamic performance and convective heat transfer for a sphere if the field is rotated with respect to the velocity vector.

1. The experiments were performed with a plasma gas-dynamic system in partially ionized nitrogen generated by a gas-discharge accelerator, in which the ionization was provided by

Dnepropetrovsk. Translated from Zhurnal Prikladnoi Mekhaniki i Tekhnicheskoi Fiziki, No. 1, pp. 15-19, January-February, 1991. Original article submitted May 15, 1989; revision submitted July 10, 1989.

