ION EXTRACTION SYSTEM STUDY

MASTER

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FINAL REPORT

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PREFACE

This document reports the results of an eight-month study contract (April through December 1976) issued by Lawrence Livermore Laboratory, and carried out under the technical direction of R. S. Lowder and M. A. Lutz. The program was devoted to the study of plasma-extraction processes for laser isotope separation of uranium. A simulation technology was developed, using mercury instead of uranium, and analytical and experimental studies of plasma-extraction processes were performed in the context of mercury simulation.

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SECTION 1

INTRODUCTION AND SUMMARY

Separation of uranium isotopes by selective laser photoionization and extraction of the photoplasma from the neutral background represents an important, energetically advantageous approach to providing enriched-uranium fuel for fission reactors. The laser separation process has the advantage that the distinction between different isotopic masses is made on a quantum level (viz., the isotopic shift in energy levels), so that only the desired mass-235 uranium atoms (^{235}U) are affected by the narrowband laser energy that is input to the system. By contrast, conventional gaseous-diffusion and centrifuge-type isotope separation plants make use of the inertial difference in isotopic masses as the distinguishing parameter. Separation by these methods typically requires from 0.45 MeV (centrifuge) to 3 MeV (diffusion) per ^{235}U atom to achieve enrichment to a reactor-grade product. This separation energy represents a significant fraction of the 200 MeV/atom fission energy of ^{235}U . In the process of laser isotope separation, the only significant inputs of energy to the particle are the heat of vaporization (4.8 eV) and the ionization energy (6 eV). The latter energy input is provided by a laser system which may be highly inefficient; if, however, the efficiency were as poor as 0.01%, only 0.06 MeV per atom would be required to achieve separation. This is an important improvement over previous technology.

Once the natural uranium has been vaporized and the mass-235 component ionized by a laser system, the resulting ^{235}U plasma must be efficiently extracted from the stream of natural uranium vapor. It is this extraction problem that has been the subject of this research program. Not only must a mechanism be identified that extracts the ^{235}U plasma without large inputs of energy, but the process must achieve a high degree of effectiveness in extraction: for every ^{235}U atom that is ionized by a 0.01% efficient laser system and does not reach the product collector, 60 keV of energy is lost.

The principal processes in the laser isotope separation process are diagrammed in Figure 1-1. Of crucial importance is the process

Figure 1-1. Block diagram of processes involved in the separation of isotopes by selective photoionization. Super-script arrows denote accelerated species.

labelled "charge exchange," in which $^{235}\text{U}^+$ ions are lost: $^{235}\text{U}^+ + ^{238}\text{U}^0 \rightarrow ^{235}\text{U}^0 + ^{238}\text{U}^+$. This loss mechanism, which is characterized by a large cross section, imposes stringent requirements on the distance through neutral vapor that ^{235}U plasma can be transported without significant losses and on the time available for extraction of that plasma. These requirements impact heavily on the choice and mode of realization of a practical plasma-extraction process.

The goal of the present program has been the identification and experimental verification (by simulation) of an effective, practical extraction mechanism which is consistent with the constraints imposed by other systems in an isotope separator, and which could be scaled to pilot-plant dimensions. This goal has been pursued by the following endeavors:

- A program of analytical study has been conducted to identify and select suitable plasma-extraction processes. These processes have been considered for their relative efficiency, ease of implementation, and compatibility with existing oven and laser system requirements. A particular approach, called "JxB-extraction" was selected as the chosen extraction method.
- An experimental facility has been constructed which permits simulation of plasma extraction processes relevant to laser-isotope separation using ^{200}Hg (mercury) in place of ^{238}U and ^{235}U . Photoionization is accomplished with a vacuum-spark ultraviolet photon source. This source is not isotope-selective; the experimental effort is directed solely toward studies of plasma extraction without regard for isotopic identity.
- A complete characterization of the properties of the source-plasmas and apparatus of the simulation facility has been conducted. This characterization has resulted in a self-consistent and repeatable data base by which plasma-extraction experiments can be gauged.
- An experimental investigation of the JxB extractor has been performed. Data acquired in this study indicate that, at the force level employed, JxB extraction is overshadowed by rapid ambipolar expansion of the Hg photoplasma. This expansion becomes directional in the presence of even a weak magnetic field.

The plasma-extraction capability of the JxB extractor has not been demonstrated in practice, and the development of this method must be left to future investigations. Data acquired during this program indicate, however, that ambipolar expansion of the photoplasma in a magnetic field may represent a plasma-extraction technology which is more than adequate for the requirements of uranium isotope separation. This program has not demonstrated a plasma extraction mechanism which is ready for isotope separation service; however, a valuable simulation technology has been established, and data have been collected which, with further research, could establish design parameters for a practical device.

SECTION 2

ION SEPARATION ANALYSIS

In this section, brief comparative analyses are presented which identify types of ion-extraction mechanisms that are suitable for application in isotope separation. Since any extraction method must be considered within the constraints of the geometry imposed by the laser and oven systems, this section begins with an identification of those constraints.

A. GENERAL CONDITIONS AND GEOMETRY FOR ISOTOPE SEPARATION

Uranium vapor, originating from an oven, streams into a laser interaction region. The vapor contains a low level of residual ionization (~0.1%) resulting from the vaporization process.

A weakly ionized plasma of the desired isotope (^{235}U) is created at the ion-extraction site by a pulse of laser radiation. This plasma is in addition to the pre-existent background plasma, which is predominantly ^{238}U . Table 2-1 gives order-of-magnitude values for the densities and temperature involved.

Table 2-1. Plasma Parameters at the Ion-Extraction Site
Immediately after Laser-firing

$$^{235}\text{U}^+ \text{ density, } n_{+,235} = 10^{17} \text{ m}^{-3}$$

$$^{238}\text{U}^+ \text{ density, } n_{+,238} = 10^{17} \text{ m}^{-3}$$

$$^{238}\text{U}^0 \text{ density, } n_{0,238} = 10^{20} \text{ m}^{-3}$$

$$^{238}\text{U}^0 \text{ temperature } T_0 = 2400^\circ\text{K}$$

1. Assay

The goal of the laser-isotope-separation approach is removal of the ^{235}U plasma from the vapor stream to a product collector with minimal loss of ^{235}U and minimum contamination of the product with natural uranium. Two natural sources of contamination exist that reduce the assay (fractional concentration of the mass 235 isotope, $\epsilon = \left[^{235}\text{U}\right] / \left(\left[^{235}\text{U}\right] + \left[^{238}\text{U}\right]\right)$ of the product. First, the background natural-uranium plasma, which is produced in the vaporization process, will be extracted along with the laser-ionized ^{235}U plasma (limiting the assay to about $\epsilon \approx 0.5$, unless this contamination is eliminated by separate processes). A second source of contamination can be found in the arrival at the product collector of natural uranium that has been scattered out of the vapor stream by thermal collisions. This source of contamination can be significantly lessened if an active ion-extraction process is employed that can deflect the laser-ionized plasma in an upstream direction. A product collector that is shielded from particles arriving from other directions can then be employed to minimize thermal-neutral contamination. The design and orientation of product collectors is beyond the scope of this report; this aspect of the total isotope separation project is considered here only in that it shows the desirability of plasma-extraction processes which are highly directional and are capable of upstream deflection of the $^{235}\text{U}^+$ ions.

In the following discussion, attention is directed toward the requirements for the extraction of $^{235}\text{U}^+$ ions from the laser-interaction region. It is assumed that some type of ion extraction process is used to accelerate $^{235}\text{U}^+$ ions out of the laser-interaction region toward a product collector. Two collisional processes which reduce product assay are examined: charge-exchange collisions and scattering of accelerated ions by momentum-exchange collisions. Requirements are calculated for the magnitude and duration of extracting forces and for the maximum dimensions of the natural-uranium vapor streamlet which can be efficiently employed.

2. Charge-Exchange Collisions

Decay of the $^{235}\text{U}^+$ plasma begins immediately after laser firing: the dominant process involved in this decay is charge-exchange with $^{238}\text{U}^0$ atoms (radiative and three-body recombination are much less probable): $^{238}\text{U}^0 + ^{235}\text{U}^+ \rightarrow ^{238}\text{U}^+ + ^{235}\text{U}^0$. Any ion-extraction process must act very rapidly so that extraction forces act on the $^{235}\text{U}^+$ plasma before it decays by charge-exchange collisions. To achieve high assay, the extraction forces must be of limited spatial and temporal duration, otherwise $^{238}\text{U}^+$ ions formed by charge exchange from the accelerated $^{235}\text{U}^+$ ions will also be accelerated toward the product collector, reducing the assay. It is therefore required that any extraction mechanism be a pulsed, short acting one.

The time-constants and characteristic distances involved in the charge-exchange process can be estimated as follows: The charge-exchange cross-section for uranium is given by $^{2-1} \sigma_{cx, U} \approx 2 \times 10^{-18} \text{ m}^{-2}$. If the $^{235}\text{U}^+$ ions had only thermal motion (i.e., were not accelerated), they would charge-exchange with the background neutrals at a rate

$$\dot{n}_{+, 235} = -n_{+, 235} n_{0, 238} \langle \sigma_{cx, U} v \rangle, \quad (1)$$

where v is the relative particle velocity and the brackets denote averaging over a Maxwellian distribution function. Using $\sigma_{cx, U}$ as given above and values from Table 2-1 in Eq. (1) gives an e-folding time of about $\tau_{cx, U} \approx 10 \mu\text{s}$ for $n_{+, 235}$. Any successful ion-extraction process must act in a time much shorter than this value. Note also that in a time $\tau_{cx, U}$, the $^{235}\text{U}^+$ ions are convected downstream a distance $\tau_{cx, U} \cdot u \approx 10 \text{ mm}$, thus imposing restrictions on product-collector design.

The charge-exchange loss process takes a somewhat different form during the ion-extraction process: if a velocity v has been imparted to the ions which is much greater than the thermal velocity v_{th} , the ions can be considered to be moving in a monoenergetic fashion through a background of stationary $^{238}\text{U}^0$ targets. The loss rate is then

$$\dot{n}_{+, 235} = -n_{+, 235} n_{0, 238} \sigma_{cx, U} v, \quad (2)$$

where v is the result of an (assumed uniform) acceleration α , $v = \dot{z} = \alpha t$. This gives

$$\frac{\dot{n}_{+,235}}{n_{+,235}} = e^{-n_{o,238} \sigma_{cx,U} \alpha t^2 / 2} \quad (3)$$

or

$$\frac{\dot{n}_{+,235}}{n_{+,235}} = e^{-n_{o,238} \sigma_{cx,U} z^2}, \quad (4)$$

where $z = \alpha t^2 / 2$ is a coordinate transverse to the vapor flow direction. Using the above values in Eq. (4) gives an e-folding distance of $\lambda_{cx,U} \approx 5$ mm. Therefore, to achieve ion extraction without substantial loss of the desired $^{235}\text{U}^+$ ions, the extraction forces must act in a time $\Delta t \ll \tau_{cx,U} \approx 10 \mu\text{s}$, and must act in such a manner that the acceleration acts over a distance $\Delta z \ll \lambda_{cx,U} \approx 5$ mm. The magnitude of the force density required is then $\mathcal{F} \gg n_{+,235} \cdot 5 \text{ mm} (10 \mu\text{s})^{-2} = 2 \text{ Nm}^{-3}$.

3. Momentum Exchange Collisions

Scattering of $^{235}\text{U}^+$ ions and charge-exchanged neutrals also occurs after the acceleration process as a result of momentum-exchange collisions with $^{238}\text{U}^0$ neutrals. The cross-section for momentum-exchange collisions is given in Ref. 2-2 as $\sigma_{mx,U} = 2 \times 10^{-19} \text{ m}^2$; therefore, the mean free path for scattering is $\lambda_{mx,U} \approx 50$ mm. It is therefore clear that, after acceleration by the ion-extraction process, ^{235}U particles (either ions or neutrals that have been formed by charge-exchange interactions with $^{238}\text{U}^0$) will still be lost if they are transported over large distances through neutral uranium vapor.

The maximum streamlet thickness that can be efficiently used in an isotope separation process is determined by the shorter of the scale lengths for plasma decay given above; this is the charge-exchange length, $\lambda_{cx,U} \approx 5$ mm. Streamlet thicknesses of this dimension or smaller must be used to avoid serious loss of $^{235}\text{U}^+$.

B. ION EXTRACTION PROCESSES (GENERAL)

On the time scales of interest in ion extraction processes for isotope separation, only two basic types of forces are available: those derived from electric and magnetic fields. This program has been concerned with two specific force mechanisms in a plasma. The first, which employs a static magnetic field with a strong spatial gradient, is referred to as the " $\mu\nabla B$ " mechanism after the force $\underline{F} = (\underline{\mu} \cdot \nabla) \underline{B}$ which acts on adiabatic charged particles in a spatial magnetic field gradient. This mechanism is representative of the forces which are available from static magnetic fields. The second mechanism considered, called the " $\underline{J} \times \underline{B}$ " mechanism, results from the interaction between induced currents in a plasma and the time-varying magnetic field that induces the currents. The $\underline{J} \times \underline{B}$ mechanism represents the class of electromagnetic (EM) forces in which the wavelength is much longer than the dimensions of the plasma. Two other classes of mechanisms which are possible — those employing short wavelength electromagnetic fields and those using purely electrostatic fields — were not studied under this program. A summary of the classes of forces available is presented in Table 2-2.

Before considering the $\mu\nabla B$ and $\underline{J} \times \underline{B}$ extraction mechanisms identified above, a discussion is given of a natural electrostatic extraction mechanism: ambipolar expansion of the photoplasma column.

Table 2-2. Classes of Force Available for Ion-Extraction

Field Type	Time Variation		
	Static	EM Long Wave	EM Short Wave
Electric	Ambipolar expansion; ion extraction from plasma boundary	$\underline{J} \times \underline{B}$; electrostatic plasma instabilities; ion heating	Electron heating; radiation pressure (classical), radiation reaction (quantum-mechanical)
Magnetic	$(\underline{\mu} \cdot \nabla) \underline{B}$, and other forces derivable from $q \underline{v} \times \underline{B}$ force		

C. AMBIPOLAR EXPANSION

The laser-generated $^{235}\text{U}^+$ plasma expands outward from its point of origin at a rate determined by time-dependent ambipolar diffusion. When charged-particle inertial forces and momentum transfer due to collisions are taken into account, the expansion process is described by the linearized second-order differential equation for the plasma number density n as a function of time and position:^{2-3, 2-4}

$$\nabla^2 n = \frac{1}{u_a^2} \frac{\partial^2 n}{\partial t^2} + \frac{1}{D_a} \frac{\partial n}{\partial t}, \quad (5)$$

where in the limit $T_e \gg T_i$,

$$u_a^2 \approx \frac{\gamma k T_e}{m_i}$$

and

$$D_a = \frac{k T_e}{m_i} \tau_i.$$

Equation (5) has been solved in one dimension by Shimony and Cahn,²⁻³ with the initial conditions

$$n(z, 0) = \begin{cases} n_0, & -r_0 \leq z \leq r_0 \\ 0, & \text{otherwise} \end{cases} \quad (6)$$

$$\left. \frac{\partial n(z, t)}{\partial t} \right|_{t=0} = 0. \quad (7)$$

The one-dimensional equation, subject to the above initial conditions, is an appropriate approximation for a situation in which a

transverse magnetic field is imposed in the region of laser interaction. In this case, diffusion parallel to the vapor flow direction is much slower than that transverse to the vapor flow. If the magnetic field is of sufficient strength, the plasma essentially undergoes free one-dimensional ambipolar expansion toward product collectors which could, for example, be placed a few mm outside the vapor stream. The fraction η of the total initial ^{235}U photoplasma* which would be collected before being lost by charge-exchange can be roughly estimated by integrating that part of the solution $n(z, t)$ to Eq. (5) that has left the vapor stream by $t < \tau_{\text{cx}, \text{U}}$:

$$\eta = 2 \frac{1}{n_0} \int_{z_0}^{\infty} n(z, \tau_{\text{cx}, \text{U}}) dz, \quad (8)$$

where $z = \pm z_0$ is the initial photoplasma boundary, and the factor 2 in Eq. (8) accounts for expansion in both the $\pm z$ directions.

The general solution to Eq. (5), subject to the boundary conditions given in Eqs. (6) and (7), is given in Ref. 2-3; the solution which is appropriate for the present conditions is described as "wave-dominated" (defined by $u_a z_0 / D_a \ll 1$). A good approximation to the behavior of wave-dominated plasma expansion, according to Ref. 2-3, is to regard the density profile as being of uniform amplitude, with the boundaries expanding at a rate $z = \pm u_a$ (see Figure 2-1). That is,

*Note that the extraction efficiency $\eta = ^{235}\text{U collected} / ^{235}\text{U ionized}$ is not directly related to assay, ϵ . Presumably, any plasma-extraction scheme extracts $^{238}\text{U}^+$ with the same efficiency η as it does $^{235}\text{U}^+$. The assay is affected by the duration of the extraction impulse compared to $\lambda_{\text{cx}, \text{U}}$ and by the volumes over which the extraction forces act (if regions of the vapor stream outside the laser-interaction region are subject to extraction forces, assay will be reduced). The importance of the efficiency η is that it appears as a direct factor in the overall energy efficiency of the isotope-separation process.

TELEGRAPHER'S EQUATION $\nabla^2 n = \frac{1}{u_a^2} \frac{\partial^2 n}{\partial t^2} + \frac{1}{D_a} \frac{\partial n}{\partial t}$

WHERE $u_a^2 \equiv \gamma \frac{kT_e}{m_i}$, $D_a = \frac{kT_e}{m_i} \tau_i$

APPROPRIATE ONE-DIMENSIONAL WAVE SOLUTION ($t < \tau_i$):

Figure 2-1. Dynamics of plasma collection resulting from ambipolar diffusion: (a) initial photoplasma; (b) after start of expansion $\Delta t < \tau_{cx,U}$; (c) late in expansion cycle, $\Delta t > \tau_{cx,U}$.

$$n(z, t) \approx \begin{cases} \frac{n_0 z_0}{z_0 + u_a t} & , \quad -z_0 - u_a t \leq z \leq z_0 + u_a t \\ 0 & , \quad |z| > z_0 + u_a t \end{cases} \quad (9)$$

Using this approximation in Eq. (9) results in a value of the fractional collection:

$$\eta = \frac{u_a \tau_{cx, U}}{z_0 + u_a \tau_{cx, U}} \quad (10)$$

Using

$$T_e = 1 \text{ eV}$$

$$z_0 = 5 \text{ mm}$$

$$\tau_{cx, U} = 5 \mu\text{s}$$

yields a value $\eta = 0.53$ (a more pessimistic value of $\tau_{cx, U} = 5 \mu\text{s}$ has been used in the calculation to allow for uncertainties in the calculated value of the charge-exchange cross-section). If the electron temperature were higher (with electron heating, for example), say $T_e = 5 \text{ eV}$, a better fraction $\eta = 0.71$ results.

These optimistic results are dimmed somewhat when one considers the requirements on the transverse magnetic field which limits the expansion to one dimension. To achieve the desired directionality requires that $u_a z_0 / D_\perp \gg 1$, where $D_\perp$ is the diffusion coefficient for ambipolar plasma expansion transverse to the magnetic field. High directionality requires $D_\perp \ll u_a z_0 = 5 \text{ m}^2 \text{ s}^{-1}$, at $T_e = 1 \text{ eV}$. Calculated expansion rates of plasma along magnetic field lines and across the field are presented in Table 2-3. These data indicate that (for modest $T_e \sim 1 \text{ eV}$) magnetic fields $B \approx 50 \text{ mT}$ are required to assure strongly directional expansion. For weaker fields ($B \approx 10 \text{ mT}$), the

Table 2-3. Ambipolar Diffusion Rates

Electron Temperature T_e , eV	Collection Efficiency, η (Calculated from Eq. 2-10)	Transport Velocity Parallel to $\underline{B}$, u_a , ms^{-1}	Transport Velocity ^a Normal to $\underline{B}$, $D_{\perp} \nabla n/n$, ms^{-1}	Type of Cross-Field Diffusion	Magnetic Field Strength, $ \underline{B} $, mT
0.2	0.33	495	220	Classical	10
			250	Bohm	10
			50	Bohm	50
1.0	0.53	1110	913	Classical	10
			1110	Bohm	10
			250	Bohm	50
10.0	0.78	3500	3500	Classical	10
			3500	Bohm	10
			2500	Bohm	50
^a $\nabla n/n$ is assumed to be 200 m^{-1} .					

parallel- and cross-field expansion velocities are calculated to be of comparable magnitude. This does not necessarily mean that expansion in a ~ 10 mT field is nondirectional. Rather, it implies that a more thorough transport analysis (including, for example, neoclassical effects and boundary perturbations) is required. Evaluation of the precise level of magnetic field strength required for one-dimensional directionality is left to experiment (cf Section 3).

D. $(\underline{\mu} \cdot \underline{\nabla})\underline{B}$ EXTRACTION

The $\underline{\mu} \cdot \underline{\nabla} B$ method of plasma acceleration is based on the formation of the $^{235}\text{U}^+$ photoplasma in a steady magnetic field having a strong spatial gradient. Under the influence of forces arising from the magnetic field gradient, electrons undergo acceleration toward the region of lower magnetic field strength. Space-charge buildup resulting from the motion of the electrons creates an electric field which accelerates ions in the direction of the electron motion, so that the entire ^{235}U photoplasma experiences a bulk force directed down the magnetic field gradient.

Electrons exposed to a spatially varying magnetic field $\underline{B}$ experience a force^{2-5, 2-6}

$$\underline{F} = q(\underline{E} + \underline{v} \times \underline{B}). \quad (11)$$

If the condition for adiabaticity

$$\frac{1}{B} \frac{dB}{dt} \ll \frac{\omega_{ce}}{2\pi},$$

where $\omega_{ce} = eB/m_e$ is the electron gyrofrequency, and $\underline{\mu}$ is the electron orbital magnetic moment, is satisfied along an electron trajectory, the magnetic field in Eq. (11) can be expanded in a Taylor's Series about the guiding center $\underline{r}_0$:

$$\underline{B}(\underline{r}) = \underline{B}(\underline{r}_0) + (\underline{v} \cdot \underline{\nabla}_0) \underline{B} + \frac{1}{2!} (\underline{v} \cdot \underline{\nabla})^2 \underline{B} + \dots$$

Identifying the gyroradius $r_c = mv_{\perp}/eB$ and the magnetic moment

$$\mu \equiv \frac{mv_{\perp}^2}{2B},$$

allows Eq. (11) to be rewritten (see geometry in Figure 2-2):

$$\underline{F} = q(\underline{E} + \underline{v} \times \underline{B}_0) - (\underline{\mu} \cdot \underline{\nabla}) \underline{B}$$

$$\underline{F}_{\mu \nabla B} = -(\underline{\mu} \cdot \underline{\nabla}) \underline{B} = -\frac{1}{2} mv_{\perp}^2 \frac{\nabla B}{B} \hat{z} = -\frac{1}{2} \frac{r_c}{L_{\perp}} v_{\perp} e B_z \hat{z}, \quad (12)$$

where

$$L_{\perp} = \left(\frac{\nabla B}{B} \right)^{-1},$$

" $\perp$ " denotes perpendicular to $\underline{B}$, and " ∇B " means $(\hat{z} \cdot \underline{\nabla}) \underline{B} \cdot \hat{z}$. The force given in Eq. (12) acts on electrons. As the electrons are accelerated away from the (nonadiabatic) ions, a space-charge separation is produced, and the resulting electric field produces a force on plasma ions of the same magnitude as that given by Eq. (12).*

*Solving the two-particle equations of motion shows that the electrons behave as if they had ionic mass. The separation between electron and ion charge sheets undergoes an oscillation at the electron plasma frequency of magnitude $\delta z \sim \mu \nabla B / m_e \omega_{pe}^2$.

To obtain an order-of-magnitude estimate of the forces achievable with this method, consider a mean electron perpendicular energy of 1 eV and a gradient $\nabla B/B$ of 100 m^{-1} . This gives a force density of $\mathcal{F} = 1.6 \text{ Nm}^{-3}$. With the same gradient at 15 eV electrons, $\mathcal{F} = 8 \text{ Nm}^{-3}$. These forces are order-of-magnitude consistent with the 2 Nm^{-3} minimum force requirement calculated in a previous section, provided that the electron temperature is sufficiently high.

1. Velocity Saturation

A difficulty arises in the $\mu\nabla B$ acceleration method when the electron temperature is not sufficiently high: acceleration in a static magnetic field gradient results from the perpendicular energy (cyclotron motion) of the electron being exchanged for energy directed along the field gradient. This exchange, at best, results in the conversion of the entire energy of the electron into motion directed along the field lines; once this point is reached, no further acceleration occurs. For low electron temperatures, the final velocity may not be adequate to remove the $^{235}\text{U}^+$ ions before charge-exchange occurs.

2. Electron Spinup

The problem of velocity saturation described above can be overcome to some measure by "spinning up" the electrons, i.e., replenishing the electron perpendicular energy $W_{\perp}$ through the use of a fast-rising magnetic field. Electrons are accelerated in the electric field $\underline{E}_1$ induced by the time-variation of the magnetic field B :

$$\underline{\nabla} \times \underline{E}_1 = - \frac{\partial B}{\partial t} \quad (13)$$

from which

$$\underline{E}_1 = - \frac{r_c}{2} \frac{\partial B}{\partial t} = - \frac{mv_{\perp}}{2eB} \frac{\partial B}{\partial t}, \quad (14)$$

where $\underline{E}_1$ is tangential to the electron orbit of radius r_c . Since $\underline{E}_1$ is always normal to $\underline{B}$, acceleration along the field lines does not result from the rising magnetic field. Time-variation of the magnetic field can, however, be exploited to increase or replenish the $W_\perp$ that is lost by conversion into parallel energy, $W_\parallel$. This conversion process occurs at a rate

$$\frac{\partial W_\parallel}{\partial t} = -\frac{\partial W}{\partial t} = v_\parallel F_\mu \nabla B = \frac{v_\parallel W_\perp}{B} (\nabla B)_\parallel \quad (15)$$

$$\frac{1}{W_\parallel} \frac{\partial W_\parallel}{\partial t} = \frac{W_\perp}{W_\parallel} v_\parallel \frac{(\nabla B)_\parallel}{B} = \nu_{\text{conv}} \quad (16)$$

The rate of change of magnetic field required to keep pace with this conversion rate is $1/B \partial B / \partial t \gtrsim \nu_{\text{conv}}$; for typical values of $W_\perp / W_\parallel$ and $\nabla B / B$, Eq. (16) gives values of $\nu_{\text{conv}} \sim 5 \times 10^5 \text{ s}^{-1}$. This rate is readily achievable in practice, so that, with a rising magnetic field, $\mu \nabla B$ acceleration beyond the original particle energy is possible, at a rate given by Eq. (12). A schematic representation of the electron spinup process is presented in Figure 2-3.

An additional benefit is realized through the use of a rapidly rising magnetic field. At the completion of laser firing, free electrons are created with perpendicular energy $W_\perp^0$, the magnitude of which depends on the details of the photoionization process employed. If, at this moment, the magnetic field B_0 is large enough for electron adiabaticity, then the electron magnetic moment $\mu = W_\perp^0 / B_0$ becomes fixed for the duration of the extraction process, an adiabatic invariant. B_0 can be chosen to be of minimal strength so that μ is maximized for a given $W_\perp^0$; as B is increased in time, μ remains constant, but the force $\underline{F} = (\mu \nabla) \underline{B}$ rapidly increases. This results in a larger total impulse to the plasma. Consider, as an example, $W_\perp^0 = 1 \text{ eV}$, $B_0 = 1 \text{ mT}$ and $\nabla B / B = 100 \text{ m}^{-1}$, and assume B rises to 10 mT in $1 \mu\text{s}$, remaining at that value afterwards. These parameters give $\mu = 1.6 \times 10^{-16} \text{ JT}^{-1}$ and an initial force density of $\mathcal{F} = 1.6 \text{ Nm}^{-3}$; however, the average force density after $5 \mu\text{s}$ is a much larger 14.6 Nm^{-3} , due to the larger magnitude of ∇B .

Figure 2-2. Geometry for discussion of $\mu \nabla B$ extraction.

Figure 2-3. Magnetic piston: plasma extraction by $(\mu \cdot \nabla)B$ with rising magnetic field. In the case of a rising field with gradient, B has been adjusted to maintain $W_{\perp}$ approximately constant.

Increasing the magnetic field strength to achieve larger average $\mu \nabla B$ is not without difficulties, however; the decade increase in B also produces a factor-of-ten increase in $W_{\perp}$, to 10 eV. As will be discussed later, this energy is well beyond the ionization limit, and it results in reduced product assay. The extent of electron spinup must be limited so as to restrict electron energy; the realizable force density is similarly restricted, as will be discussed in Section 2-F.

3. Collisional Losses

The $\mu \nabla B$ acceleration process is slowed by electron collisions, which isotropize the electron distribution function in the lab frame. Loss of directed electron velocity occurs at a rate $\nu_c \sim \nu_{eo}$ close to the electron-neutral collision frequency, because the electron distribution function already appears isotropic to co-moving electrons and ions. Electrons therefore experience a net force (acceleration minus viscous drag) of

$$F_{\text{tot}} = W_{\perp} \frac{\nabla B}{B} - \nu_c m_e v_{eo}. \quad (17)$$

Setting $F_{\text{tot}} = 0$ in Eq. (17) determines a maximum velocity

$$v_{\text{max}} = \frac{W_{\perp}}{m_e \nu_{eo}} \frac{\nabla B}{B}. \quad (18)$$

For reasonable values of the parameters in Eq. (18), v_{max} is extremely large; collisional isotropization does not represent a significant drag mechanism at the modest velocities required for extraction.

In summary, it has been found that the field-gradient acceleration mechanism provides a bulk force of only marginal magnitude. This force requires a magnetic field with a large spatial gradient that is also rapidly rising in time. The force saturates by collisions only at very high velocities.

E. $\underline{J \times B}$ EXTRACTOR

Electrons subject to a fast-rising magnetic field undergo $\underline{ExB}$ drifts in response to the instantaneous magnetic field $\underline{B}$ and the induced electric field $\underline{E}_1$. The value of the induced electric field is calculated from Maxwell's equation,

$$\nabla \times \underline{E}_1 = - \frac{\partial \underline{B}}{\partial t} \quad (19)$$

from which one obtains

$$\underline{E}_1 = - \hat{y} \int \frac{\partial}{\partial t} B(x, z) dx + f(z) \approx - \dot{B}_z L_2 \hat{y}. \quad (20)$$

Here L_2 is the distance from the magnetic axis, as shown in Figure 2-4. The electric field $\underline{E}_1$ establishes an electron drift $\underline{v}_D = \underline{ExB}/B^2$ (which is again due to the nonadiabaticity of the ions, which lag behind); this drift produces a charge separation and an electric field $\underline{E}_2 \hat{z}$. The field $\underline{E}_2$ is of such strength that the $\underline{E}_1 \times \underline{B}$ drift ceases; that is, the effect of the magnetic field on particle motion is exactly counteracted by the field $\underline{E}_2$. Electrons accelerated in the $\hat{y}$ direction by $\underline{E}_1$ are deflected in the $\hat{z}$ direction by the Lorentz $q \underline{v} \times \underline{B}$ force; however, $\underline{E}_2$ provides an opposite deflection of such magnitude that there is no net electron drift motion in the $\hat{z}$ direction. See Figure 2-5 for a graphical illustration of the above discussion. The electric field $\underline{E}_2$ has magnitude^{*}

$$\underline{E}_2 = v_{\perp} B_x \hat{z}, \quad (21)$$

where $v_{\perp}$ is the component of electron velocity normal to $\underline{B}$. The velocity $v_{\perp}$ results from acceleration in the electric field $\underline{E}_1$, and is limited by collisions:

^{*}Ion acceleration occurs in the electric field $\underline{E}_2$ ($\underline{Z} = e \underline{E}_2 / m_e$). Comparing this force (where $\underline{E}_2$ is given by Eq. (21)) with the $\mu \nabla B$ force given by Eq. (12) reveals that the magnitude of the $\underline{J \times B}$ force exceeds that of the $\mu \nabla B$ force by $2 L_1 / r_c \gg 1$, where B_x is assumed $\approx B_z$.

Figure 2-4. Geometry and coordinate system for discussion of $\underline{J} \times \underline{B}$ forces.

Figure 2-5. Conceptual sequence of events leading to $\mathbf{J} \times \mathbf{B}$ plasma acceleration. Electron current paths are closed in the lower half space (not shown), where $\mathbf{B}$ is directed downward.

$$\frac{dv_1}{dt} = \frac{eE_1}{m_e} - \nu_{eo} v_e, \quad (22)$$

where ν_{eo} is the collision frequency for electron-neutral collisions, and v_e is the total electron velocity, $v_e^2 = v_1^2 + kT_e/m_e$. For low initial electron temperature ($T_e \lesssim 1$ eV), Eq. (22) describes a rapidly rising v_1 . Large-angle collisions with neutrals rapidly thermalize the electron distribution function, however, producing electron heating. By the time $T_e \gtrsim 1$ eV, $v_1 \ll v_e$, and v_1 can be calculated approximately by a magnetic-field-free mobility:

$$v_1 = \frac{eE_1}{m_e \nu_{eo}} = \frac{eE_1}{m_e n_{o,238} \sigma_{eo} v_e}, \quad (23)$$

where σ_{eo} is the cross-section for electron-neutral collisions, and $v_e \approx \sqrt{2kT_e/m_e}$. Substituting $\underline{E}_1 = \dot{B}_z L_2$, and using Eq. (21) yields

$$\underline{E}_2 = \frac{e B_x \dot{B}_z L_2}{m_e n_{o,238} \sigma_{eo} v_e} \hat{z}, \quad (24)$$

from which the force density on the ions can be written:

$$\underline{F}_{JxB} = -en_{+,235} \underline{E}_{JxB} = \frac{e^2 B_x \dot{B}_z L_2 n_{+,235}}{m_e n_{o,238} \sigma_{eo} v_e}. \quad (25)$$

To estimate the magnitude of force involved here, take

$$B_x \approx B_z = 10 \text{ mT}$$

$$L_2 = 5 \text{ mm}$$

$$\dot{B}_z = 10^4 \text{ T s}^{-1}$$

$$\sigma_{eo} v_e = 8 \times 10^{-13} \text{ m}^3 \text{ s}^{-1} \quad (T_e \approx 1 \text{ eV}).$$

These values and the densities given previously yield

$$\underline{\mathcal{F}}_{\underline{JxB}} \approx 18 \text{ Nm}^{-3}.$$

This force density, $\underline{\mathcal{F}}_{\underline{JxB}}$, is an order of magnitude larger than the minimum force density of 2 Nm^{-3} calculated in an earlier section as a requirement for an isotope separation process which would function successfully in the face of a high charge-exchange rate. The $\underline{JxB}$ extraction method differs from the $(\underline{\mu} \cdot \underline{\nabla})\underline{B}$ mechanism primarily in that macroscopic particle trajectories are emphasized rather than microscopic, single-electron orbits.

F. COMPARISON OF PLASMA EXTRACTION METHODS

The preceding discussion of the three plasma-extraction processes can be summarized as follows:

- Ambipolar diffusion in a magnetic field transverse to the vapor flow direction has been shown to be adequate for isotope separation purposes provided that the electron temperature is high enough (or if rapid supplemental heating is provided).
- The $(\underline{\mu} \cdot \underline{\nabla})\underline{B}$ extraction force has been calculated to be only marginally adequate, and then only with high spatial magnetic-field gradients and rapidly rising magnetic field strength.
- Calculations of the $\underline{JxB}$ body force result in values of plasma acceleration well above the minimum required value, using readily obtainable values of B and $\underline{B}$.

All of these processes are subject to a common limitation: the electron energy must not be allowed to approach the value (6.08 eV) at which collisional ionization of $^{238}\text{U}^0$ atoms occurs; it is easily calculated, for example, that if 10% of the electron population has an energy in excess of the $^{238}\text{U}^0$ ionization potential for a 5 μs extraction time, the product assay will be reduced by a factor of three. Accordingly, it is required that the application of extraction forces (all of which function better at high electron energy) be carried out in such a manner that electron energy is restricted below the ionization limit.

Nonclassical particle transport is potentially an additional restriction on the magnitude of the extraction forces. Although the effects of nonclassical transport have been discussed only for the case of plasma extraction by one-dimensional ambipolar expansion, these effects could also act to reduce the efficacy of the $\underline{J \times B}$ and $\underline{\mu \nabla B}$ extractors. No attempt has been made to calculate neoclassical or wave-enhanced transport coefficients in the comparisons which follow, the investigation of such processes being left to experiment.

In comparing the plasma-extraction methods given above, it is necessary to assess not only the ease with which a given extraction force is produced, but also the magnitude of the force generated when the electron energy is maximal.

1. Comparison of Forces Achieved by $(\underline{\mu} \cdot \underline{\nabla})\underline{B}$ and $\underline{J \times B}$

Equations were given in earlier sections for the forces developed by the $(\underline{\mu} \cdot \underline{\nabla})\underline{B}$ and $\underline{J \times B}$ mechanisms; these are repeated below for convenience.

Extraction Forces (per particle):

$$F_{\underline{\mu \nabla B}} = \frac{1}{2} \frac{r_c}{L_1} v_{\perp} e B_z \quad (26)$$

$$F_{\underline{J \times B}} = e v_{\perp} B_x \quad (27)$$

Electric Fields:

$$E_{\underline{\mu \nabla B}} = \frac{\dot{B}_z \gamma_c}{z} \quad (28)$$

$$E_{\underline{J \times B}} = v_{\perp} B_x \quad (29)$$

Consider first the force equations, Eqs. (26) and (27). Each equation contains a perpendicular electron velocity $v_{\perp}$, which is the result of acceleration in an induced electric field. These velocities must be maintained below an upper bound $v_{\perp} < v_{\max}$, where $v_{\max}$ is the velocity at which electron-impact ionization of neutral-uranium atoms becomes unacceptably large. Assuming that $v_{\perp} = v_{\max}$ in both cases, and assuming $B_z \approx B_x$, Eqs. (26) and (27) yield

$$\left. \frac{F_{\underline{J \times B}}}{F_{\underline{\mu \cdot \nabla B}}} \right|_{v_{\perp} = v_{\max}} = \frac{2L_1}{r_c} . \quad (30)$$

That is, the force-per-ion produced by the $\underline{J \times B}$ mechanism is $2L_1/r_c$ times greater than that produced by the $(\underline{\mu \cdot \nabla})\underline{B}$ mechanism. In the ideal collisionless case, this ratio can be made arbitrarily large; however, particle collisions and scattering from turbulent electric fields can result in the perturbation of macroscopic particle trajectories, limiting the achievable value of $2L_1/r_c$.

2. Comparison of Magnetic-Field Requirements for $\underline{\mu \cdot \nabla B}$ and $\underline{J \times B}$ Mechanisms (Collisionless Approximation)

The comparison of forces carried out above has assumed that equal electron velocities $v_{\perp} = v_{\max}$ were reached in each case. If we instead consider the case in which equal forces have been achieved (again assuming $B_x \approx B_z$), the required value of $\dot{B}$ can be calculated for each case; i. e., assume

$$F_{\underline{\mu \cdot \nabla B}} = F_{\underline{J \times B}}$$

Using Eqs. (26) and (27), this requires that

$$v_{\perp, \underline{J \times B}} = \frac{r_c}{2L_1} v_{\max} . \quad (31)$$

In the $\underline{J \times B}$ mechanism, the velocity v_1 develops as a function of time from the electric field $E_1 = \dot{B}_z L_2$; the (collisionless) velocity is given by

$$v_{1, \underline{J \times B}} = \frac{-e L_2}{m_e} t \dot{B}_{z, \underline{J \times B}} \quad (32)$$

Similarly, in the $(\underline{\mu \cdot \nabla})B$ acceleration mechanism,*

$$v_{1, \underline{\mu \cdot \nabla} B} = v_{\max} = - \frac{e r_c}{2 m_e} t \dot{B}_{z, \underline{\mu \cdot \nabla} B} \quad (33)$$

From Eqs. (32) and (33), one can conclude (taking $L = L_1 \approx L_2$) that

$$\left. \frac{\dot{B}_{z, \underline{J \times B}}}{\dot{B}_{z, \underline{\mu \cdot \nabla} B}} \right|_{\text{EQUAL FORCES}} \approx \left(\frac{r_c}{2L} \right)^2 \quad (34)$$

Equation (34) reveals that the $\underline{J \times B}$ acceleration mechanism achieves a force equal to that produced by the $(\underline{\mu \cdot \nabla})B$ mechanism with a much lower value of $\dot{B}_z$. Equation (34) is valid in the collisionless regime, i.e., for times short with the electron collision time, τ_e : $t \ll \tau_e$.

3. Comparison of Magnetic Field Requirements for $\underline{\mu \cdot \nabla} B$ and $\underline{J \times B}$ Mechanisms (Collision-Dominated Case)

For room temperature electrons, τ_e is of the order of 1 μs ; however, for $T_e = 1$ eV, $\tau_e = 1/\nu_e \approx 30$ ns. In this case, the velocities given in Eqs. (32) and (33) immediately saturate at collision-dominated values. The limiting velocities can be found by using a magnetic-field-free mobility:

*Equation (33) is a low-order approximation, because $r_c = mv_1/eB$ is time-varying and has been treated as constant. The r_c in Eq. (33) may be thought of as an average gyroradius over the time of interest. Arguments based on Eq. (33) are order-of-magnitude in nature and are unaffected by this approximation.

$$v_{1, \underline{JxB}} = \frac{e E_1}{m_e v_e} = \frac{e E_1}{m_e n_o \sigma_{eo} v_e}$$

$$v_{1, \underline{JxB}} = \frac{e \dot{B}_z L_2}{m_e n_o \sigma_{eo} v_e} . \quad (35)$$

Similarly, for the $(\underline{\mu \cdot \nabla})B$ case:

$$v_{1, \underline{\mu \nabla B}} = v_{\max} = \frac{e \dot{B}_z r_c}{m_e n_o \sigma_{eo} v_e} . \quad (36)$$

Combining Eqs. (31), (35), and (36) results in a collision-dominated analog to Eq. (34):

$$\frac{\dot{B}_z, \underline{JxB}}{\dot{B}_z, \underline{\mu \nabla B}} = \left(\frac{r_c}{2L} \right)^2 . \quad (37)$$

The advantage of the $\underline{JxB}$ acceleration mechanism is equally as profound in the collision-dominated case as in the collisionless one. The comparisons described above, and the force terms and effective electric field terms for each acceleration mechanism, are summarized in Table 2-4.

Calculations performed above have established that the $\underline{JxB}$ acceleration mechanism is greatly superior to the acceleration mechanism which is dependent upon the $(\underline{\mu \cdot \nabla})B$ force. The superiority is apparent in that a higher value of ion acceleration can be achieved before the onset of electron-impact ionization of neutrals. It is also apparent in that ion acceleration is achieved with a lower value of $\dot{B}$, so that the technological requirements are more easily met for this type of acceleration.

The discussion of the previous sections establishes a clear mandate for the adoption of the $\underline{JxB}$ extraction mechanism as the chosen approach. This is not to say, of course, that ambipolar diffusion and $\underline{\mu \nabla B}$ forces will not be present, but rather that experimental geometry will be

Table 2-4. Summary Comparison of $\underline{J \times B}$ and $(\underline{\mu \cdot \nabla})\underline{B}$ Acceleration Mechanisms

Acceleration Mechanism	Effective Perpendicular Electric Field, $\frac{\oint \underline{E} \cdot d\underline{\ell}}{\oint d\underline{\ell}}$	Extraction Force on Ion, F_z	Relative Force on Ion, Assuming $v = v_{\text{max}}$	Relative B_z Requirement, Assuming Equal Extraction Forces on Ions, F_z	
				Collisionless ^a	Collisional
$(\underline{\mu \cdot \nabla})\underline{B}$	$\frac{\dot{B}_z r_c}{2}$	$\frac{ev_{\perp} B_z r_c}{2L_1}$	1	1	1
$\underline{J \times B}$	$\dot{B}_z L_2$	$ev_{\perp} B_x$	$\frac{2L_1}{r_c}$	$\left(\frac{r_c}{2L_1}\right)^2$	$\left(\frac{r_c}{2L_1}\right)^2$

^a See footnote following Eq. (32).

designed to best exploit the $\underline{J \times B}$ extraction force. Such a geometry is shown in Figures 2-6 and 2-7. A single-turn, flat coil is placed near the ^{235}U photoplasma. Current flowing out of the page in the upper part of the coil induces an oppositely directed plasma current $\underline{J}$, which interacts with vertical magnetic field lines near the coil surface, providing the $\underline{J \times B}$ extraction force (as illustrated earlier in Figure 2-5). This situation is duplicated in the lower portion of the coil, and the presence of the curved ends of the drive coil produces a plasma current consisting of closed electron paths. Closure of the electron current path is an important aspect of the design of a $\underline{J \times B}$ extractor; open paths will result in the buildup of space charge at plasma boundaries. Space-charge buildup continues until an electric field is developed which cancels the electric field induced by the time-variation of B ; when this occurs, there is no plasma current and no net force on the plasma.

A second condition is required to assure that ions as well as electrons are accelerated by the $\underline{J \times B}$ mechanism; electrons from the background thermal ^{238}U plasma must be isolated from the ^{235}U plasma. Since the $\underline{J \times B}$ force acts only on the electrons, ions are accelerated only because of the space-charge imbalance resulting from electron depletion. If background electrons can flow freely into the extraction region, the electron depletion will be neutralized, and no ion acceleration will take place. As shown in Figure 2-4, this problem is avoided by the component of the magnetic field which is normal to the vapor flow; background electrons must diffuse across the magnetic field to reach the extraction region. Even if this crossfield diffusion is non-classical (Bohm), a weak magnetic field is sufficient to ensure that ions are extracted faster than background electrons can diffuse into the extraction region.

Figure 2-6. Geometry for the $\underline{J} \times \underline{B}$ extractor.

Figure 2-7. Chosen geometry for ion extraction.

SECTION 3

ION SEPARATION EXPERIMENTS

The experimental phase of the ion-extraction program is discussed in this section. Apparatus was constructed to simulate the conditions and geometry of a uranium isotope separation device; mercury is used in place of uranium, and ionization by UV photons (from a vacuum spark source) is used to simulate laser photoionization. No attempt is made to preferentially ionize or separate a particular isotope; rather the focus of the experimental program was the establishment of a plasma-acceleration process to enable rapid removal from a vapor stream of ions formed by the photo-ionization process.

A. PLASMA EXTRACTION APPARATUS

An overall view of the experimental plasma-extraction facility is shown in Figure 3-1. A mercury plasma source (shown in Figure 3-2) supplies weakly ionized (or neutral, as desired) mercury vapor from a thermal vaporation source and reflex discharge. The mercury flow rate and discharge power are adjusted to provide a nominal

$$n_{o,Hg} = 10^{20} \text{ m}^{-3}$$

and

$$n_{+ ,Hg} = 10^{17} \text{ m}^{-3}$$

at a location approximately 50 mm below the vapor source exit nozzle (a slit nozzle, 4 mm x 40 mm). At this location a beam of photons passes through the mercury vapor, creating a cylindrical region of enhanced plasma density due to photoionization of the mercury vapor. The vapor stream is directed downward into a liquid-nitrogen

Figure 3-1. Experimental plasma-extraction apparatus.

Figure 3-2. Mercury plasma source.

filled cryo-collector. This collector freezes the mercury vapor, thus maintaining a one-directional flow of vapor. A large liquid-nitrogen-filled cryowall is also present in the vacuum vessel to pump background mercury vapor.

1. Photon Source

Photoionization of the mercury vapor is accomplished with a beam of photons emanating from a vacuum spark source,³⁻¹ shown in Figure 3-3. The source consists of a pair of hollow electrodes separated by a 57 mm long alumina cylinder. Typically, 5 to 10 kV is placed across the electrode pair, and a discharge is initiated by a spark from a trigger electrode which is coaxial with the off-ground high-voltage electrode. The trigger electrode is insulated from the high-voltage electrode by a teflon sleeve; this sleeve is eroded during the main discharge, the teflon forming a high-density plasma from which a copious UV photon flux is emitted. The photons escape through the bore of the grounded electrode, then pass through a 1 m long evacuated* collimating tube. The other end of the tube is fitted with a 4 mm diameter circular aperture, which serves as a beam-forming collimator (the resulting photon beam has a divergence half-angle of less than 4.5 mrad. A permanent magnet with field-shaping pole-pieces is placed midway down the collimating tube. This magnet creates a strong transverse magnetic field, which effectively prevents plasma formed in the photon-source discharge from entering the main vacuum vessel.

Output of the source consists of a burst of UV photons of approximately 1 μ s duration, with a 2 μ s long "tail." An oscilloscope recording of the photon output is shown in Figure 3-4, and a radial profile of photon-beam intensity is plotted in Figure 3-5. These data were collected by observing the photoelectron current from a negatively biased Langmuir probe as it was moved across the beam (in the absence of flowing mercury vapor).

*Up to 1 μ m of Ar pressure is often used in this tube; this flowing argon results in more repeatable source operation.

Figure 3-3. UV photon source for photoplasma generation.

Figure 3-4.
Photon source output, as detected by negatively biased Langmuir probe (tungsten, $0.13 \text{ mm} \times 2.5 \text{ mm}$). Ten shots were fired in this multiple exposure to show repeatability.

Figure 3-5. Radial profile of photoelectron current collected by a Langmuir probe placed in the photon beam. A beam collimator, which normally restricts the beam to a 4 mm diameter, was removed for these measurements.

The light output of the photon source is incident on a shielded metal-disc photon collector at the opposite end of the vacuum vessel. Energetic UV photons (energy $h\nu \approx 5$ eV) release photoelectrons from the collector at a rate of about 10^{-2} electrons/photon (see Figure 3-6). A calculation of the total UV photon output, $\Gamma_Y \Delta t$, in the 4 mm exit beam has been made from these data, resulting in a value $\Gamma_Y \Delta t \approx 2 \times 10^{16}$ photons/m². This value corresponds to a source efficiency of about 1% (photon energy/capacitor energy). Operation of the source is characterized by a very high degree of repeatability, permitting the recording of point-by-point Langmuir probe characteristics.

2. Diagnostics

Two principal diagnostic tools are employed to measure properties of the mercury discharge plasma: Langmuir probes and a Faraday cup. Two Langmuir probes are mounted, each consisting of a tungsten wire 0.13 mm diameter by 2.5 mm long protruding from the end of a 1 mm diameter rigid coaxial cable. Langmuir probe characteristics are obtained by biasing the probe to each of several values of DC bias voltage, firing the photon source, and plotting the probe current at a given time after source-fire as a function of bias voltage. The high shot-to-shot repeatability of the photon source and the stability of the mercury vapor/plasma source make possible the recording of a complete, consistent Langmuir probe characteristic, using separate firings of the source for each point on the curve. Each Langmuir probe characteristic provides the electron temperature, electron plasma density, and plasma potential at a particular point in space and time. The probes are movable from outside the vacuum wall via a three-axis traverse, so that probe measurements of the entire plasma-extraction region can be taken without breaking vacuum. A typical Langmuir probe characteristic is shown in Figure 3-7, where the method of reduction is indicated.

A Faraday cup has also been used to measure properties of streaming mercury plasma. This device consists of a negatively biased collecting plate (to collect ions) preceded by a secondary-electron

Figure 3-6. Photoelectron yield of tungsten as a function of photon energy. Adapted from Ref. 3-2.

Figure 3-7. Typical Langmuir probe characteristic, showing method of data reduction. Each point plotted corresponds to a separate source firing. These data were taken on the photon beam axis, 5 μs after the start of source-firing.

repelling grid, and housed in a grounded coaxial shield with grounded entrance grid. Both grids are of high transparency (82%), photoetched nickel mesh. The characteristic mesh dimension is smaller than a Debye length in the plasmas measured. The Faraday cup has been used to measure ion currents produced by the plasma acceleration process. It provides a rapid and convenient detector of moving plasma when density and temperature information is not required. In addition, the Faraday cup can be viewed as a surrogate product collector, the "product" deposition being determined from the total ion charge received.

B. SOURCE PLASMA CHARACTERIZATION

The experimental plasma-extraction apparatus described in the previous section produces a stream of flowing mercury vapor which is partially ionized by a UV photon beam from the photon source. In addition, the plasma from the reflex discharge source streams into the plasma-extraction region through a glass nozzle, in simulation of the background natural-uranium plasma leaving the uranium oven. Results are presented in this section of a program of characterization of these plasmas. The measured plasma properties are compared with expected properties based on published cross-section data, etc., to assure that the data are consistent and meaningful. To further confirm consistency of the data, these measurements were repeated using argon in place of mercury, although all plasma-extraction experiments were conducted with mercury only.

1. Photoplasma Measurements with Mercury

Characteristics of the photoplasma produced in a flowing mercury stream have been documented over a wide range of mercury flowrates. Characteristics like that in Figure 3-7 have been recorded for several mercury flowrates; the results of these measurements are shown in Figure 3-8, where on-axis plasma density is plotted against the density of the flowing mercury vapor.

Figure 3-8. Variation of the density of mercury and argon plasma with neutral mercury and argon density, respectively. These data were taken on the photon beam axis $5 \mu\text{s}$ after the start of source firing. The neutral densities were calculated from flowrates and equilibrium temperatures.

In Figure 3-9 a radial profile of the plasma density is shown. These data were taken 5 μ s after the start of source firing. The distribution of plasma density in this figure suggests that considerable expansion of the plasma column has taken place since source firing. Integrating the area under the curve in Figure 3-9 gives the total number of particles:

$$N = \int_0^{\infty} n_{Hg^+}(r) 2\pi r dr \approx 1.1 \times 10^{11} m^{-1}.$$

if these particles were contained within the original 4 mm beam diameter, an "initialized" plasma density

$$n_{Hg^+}^0 = 8.75 \times 10^{15} m^{-3}$$

would result, where

$$n_{Hg^+}^0 = N / \left(\frac{\pi}{4} (4 \text{ mm})^2 \right).$$

This "initialized" density is the initial density which would have been produced if the photon burst had been of infinitely short duration.

2. Photoplasma Measurements with Argon

The lower curve in Figure 3-8 gives the variation of argon plasma density as a function of the density of flowing argon. As in the case with mercury, the data were taken on the photon beam axis 5 μ s after source firing. The yield of argon plasma from a given neutral particle density and UV-photon flux is lower than that for mercury by almost a factor of two. Beam-profile data were not taken for the argon plasma,

Figure 3-9. Radial profile of the density of mercury plasma, measured by Langmuir probe 5 μ s after the start of source firing.

but if the profile is assumed to be of the same shape as that measured with mercury, an initialized density

$$n_{Ar^+}^0 = \frac{\int_0^\infty n_{Ar^+}(r) 2\pi r dr}{\frac{\pi}{4} (4 \text{ mm})^2} = 4.75 \times 10^{15} \text{ m}^{-3}$$

results.

3. Reprise of Photoplasma Measurements

The initialized density n_i^0 of a plasma created by photoionization (before plasma expansion has taken place) is given by

$$n_i^0 = \langle \Gamma_Y(\nu) \sigma_{oi}(\nu) \rangle n_o \Delta t, \quad (38)$$

where $\Gamma_Y(\nu)$ is the flux of incident photons of frequency ν , $\sigma_{oi}(\nu)$ is the photoionization cross-section at frequency ν , n_o is the neutral-particle density, Δt is the duration of the photon burst, and the brackets denote averaging over the photon frequency distribution. The plasma density n_i^0 may be estimated from Eq. (38) for the present experimental conditions by employing the value given earlier: $\Gamma_Y \Delta t = 2 \times 10^{16} \text{ m}^{-2}$. This flux corresponds to the energy range above the work function of the metal surface, $h\nu \gtrsim 5 \text{ eV}$. Photoionization cross-sections for mercury and argon, taken from Refs. 3-3, 3-4, and 3-5, are shown in Figure 3-10. In the energy range of interest, the following approximate values are appropriate:

$$\text{Argon: } \sigma_{oiAr} = 3 \times 10^{-21} \text{ m}^2$$

$$\text{Mercury: } \sigma_{oiHg} = 6 \times 10^{-21} \text{ m}^2.$$

Figure 3-10.
Photoionization cross sections of mercury and argon. Top: mercury, relative cross section adapted from Ref. 3-4, using a value for the absolute cross section at threshold from Ref. 3-5. Bottom: argon, adapted from Ref. 3-3.

Using these cross-section values in Eq. (38) gives calculated plasma densities which are compared with observed initialized values in Table 3-1.

Table 3-1. Comparison of Observed and Calculated Initialized Densities n_i^0

Gas	Initialized Plasma Density, n_i^0	
	Calculated, $\Gamma_Y \Delta t n_o \sigma_{oi}$	Measured, $\frac{\int_0^\infty n_i(r) 2\pi r dr}{\frac{\pi (4 \text{ mm})^2}{4}}$
Argon ^a	$6 \times 10^{15} \text{ m}^{-3}$	$4.8 \times 10^{15} \text{ m}^{-3}$
Mercury	$1.2 \times 10^{16} \text{ m}^{-3}$	$8.7 \times 10^{15} \text{ m}^{-3}$
^a Argon radial profile assumed to be the same as observed with mercury.		

The data in Table 3-1 show that the observed rates of photo-plasma generation are consistent with those calculated from the observed photon flux. These rates are significantly below the value $n_{Ar^+}^0 = 10^{17} \text{ m}^{-3}$ reported in Ref. 3-1; the discrepancy between that datum and present results may be attributed to a lower output from the photon source. The plasma-extraction experiments which are described in the next section were conducted with a mercury photoplasma density somewhat below the goal value of $n_i^0 = 10^{17} \text{ m}^{-3}$. Results are also presented for extraction of the streaming discharge plasma, at higher density, $n_i \geq 10^{17} \text{ m}^{-3}$.

C. ION-EXTRACTION EXPERIMENTS

Experimental results are presented in this section for the extraction of Hg^+ ions from flowing mercury vapor which has been photoionized. The apparatus is shown in Figure 3-11.

In the following discussion, results are first presented which characterize the free expansion of the photoplasma column in the absence of the drive coil magnetic field. A second section describes the drive coil itself and expected extraction force, and finally, measurements which characterize the effect of the coil are presented.

1. Faraday-Cup Measurements of Ambipolar Expansion

In the absence of the drive coil magnetic field, the photoplasma column expands freely in all directions; this expanding plasma is readily collected by the Faraday cup (shown in Figure 3-11). As the Faraday cup is moved away from the plasma column (increasing z), the total charge $C = \int Idt$ collected by the cup diminishes as the cup subtends a smaller and smaller solid angle to the photon beam axis. In Figure 3-12, the dependence on z of the charge Q collected by the Faraday cup is shown for the case of free expansion ($B = 0$). A curve showing the expected charge (calculated from the collection geometry and fit to the lowest $-z$ point) is shown for comparison. The observed dependence agrees well with the calculated dependence except at large z , where the observed charge collection falls off faster with z than predicted. The parameter required in fitting the predicted-charge curve to the first data point is the initialized plasma density n_i^0 . The value required was $n_i^0 = 2 \times 10^{15} \text{ m}^{-3}$; this density is in good agreement with that predicted on the basis of measured neutral mercury flowrate and with the observed ionizing power of a photon pulse.

The stronger than predicted dependence of Q on z , which is apparent in Figure 3-12, may be due to charge losses. The resonant charge-exchange length in mercury at this density is of the order of 10 to 20 mm;³⁻⁶ ions which undergo charge exchange may be preferentially convected away in the flowing mercury vapor.

Figure 3-11. Arrangement of apparatus for ion-extraction experiments. A pair of Langmuir probes, movable over the entire extraction region is not shown. A coordinate system with origin on the photon beam axis is indicated.

Figure 3-12. Measured dependence of the charge collected by the Faraday cup on distance from the photon beam axis. The solid curve represents the charge expected on the basis of (1) an initialized density of $n_i = 2 \times 10^{15} \text{ m}^{-3}$ ($Q_{total} = 149 \text{ pC}$) and (2) the collection geometry. In the equation, D = diameter of Faraday cup, and L = photoplasma column length.

These data establish the consistency with expectations of charge measurements with the Faraday cup, providing a useful benchmark for comparison with measurements of plasma behavior under the influence of the drive coil magnetic field.

2. Drive Coil Magnetic Field

Measurements of the magnetic field produced by the drive coil (shown in Figure 3-13) are presented in Figure 3-14; these values were measured with a Gaussmeter under DC current conditions, and scaled to the 1 kA current used in the experiments.* This scaling applies a few μ s after coil firing, when the current pulse has fully penetrated the drive coil. The $\underline{J \times B}$ force to be expected from the coil can be readily calculated from these data by using Eqs. (21) and (23). At a coil current $I = 800$ A and for $z = 10$ mm and $x = L_2 = 6$ mm:

$$B_x = 2.8 \text{ mT}$$

$$B_z = 9.2 \text{ mT}$$

$$\dot{B}_z \approx 400 \text{ Ts}^{-1}.$$

Using Eq. (23), one finds $v_1 = 4.2 \times 10^3 \text{ ms}^{-1}$; from Eq. (21), one finds a force per particle of $F = 1.9 \times 10^{-18}$ N. The acceleration of an ion of mass 200 (Hg^+) is therefore $\ddot{z} = 5.7 \times 10^6 \text{ ms}^{-2}$. This acceleration if continued at a constant value over a distance of 10 mm would result in an ion velocity in excess of 330 ms^{-1} . The force produced by this coil, which corresponds to a force density of $\mathcal{F} = 0.2 \text{ Nm}^{-3}$ at $n_e = 10^{17} \text{ m}^{-3}$, extends only over a limited region of space. The small dimensions of the coil are dictated by the requirement that electron current paths be closed within the boundaries of the plasma. In the present experiments, the achievable force density $\mathcal{F}_{\underline{J \times B}}$ is limited by component parameters in the drive coil supply circuit.

*Measurements of the EMF produced by the pulse coil were also carried out, using a 10 mm diameter pickup loop. At the location $z = 10$ mm, $x = 6$ mm considered above, the peak induced voltage was found to be 1 V per turn. These data cannot be directly compared with those given above, however, because they correspond to a time soon after coil firing when the current pulse had not yet fully penetrated the drive coil.

Figure 3-13. Magnetic piston drive coil and firing circuit.

Figure 3-14.
Variation with position z of the x - and z - components of the drive coil magnetic field, B_x and B_z , at a coil current of 1 kA. Data have been scaled from dc measurements at 10 A.

3. Effect of The Magnetic Field on Photoplasma Expansion

Figures 3-15 and 3-16 display the time evolution of the density of plasma detected by Langmuir probes placed outside the initial boundaries of the photoplasma column. These data were taken without the flowing background plasma which is normally provided by the reflex discharge source; instead, an electrically-emissive tungsten filament was suspended in the photoplasma beam. This filament was found necessary to achieve repeatable results when operating without the DC background plasma; it is felt that, in the absence of the filament (or other electron source), stray electric fields arising from plasma-wall interactions suspended in the photoplasma beam. This filament was found necessary to secure repeatable results when operating without the DC background plasma; it is felt that in the absence of the filament (or other electron source) stray electric fields arising from plasma-wall interactions perturb plasma expansion. These perturbations are minimized by providing a stable source of ground-potential electrons. The data in Figure 3-16 were acquired with a Langmuir probe located 10 mm away from the optical centerline, in the direction of expected plasma motion (transverse to the direction of vapor flow). The data points indicated by circles, taken without firing the drive coil, represent free expansion of the photo-plasma column. The data points indicated by triangles were taken with the drive coil operational; increased plasma density late in time is evident in these data as a result of plasma acceleration by the magnetic piston. For all data presented here, the drive coil has been triggered from the photon burst itself (detected by a remote Langmuir probe), and the coil current rises from zero to 1 kA in 14 μ s. Accordingly, in the present experiments, the dominant effect of the magnetic piston appears late in time ($t \approx 4 \mu$ s), after generation of the photoplasma is complete ($\sim 2 \mu$ sec).

In Figure 3-15, another pair of time-resolved plasma density curves is shown, this time taken with a Langmuir probe placed immediately below the photoplasma column. A significant reduction in plasma density late in time is evident in the data in which the magnetic piston has been fired.

Comparison of Figures 3-15 and 3-16 shows that the magnetic field effectively increases plasma flow in the direction transverse to the vapor flow, at the expense of that parallel to the vapor flow.

Figure 3-15. Electron plasma density measured by a Langmuir probe located 10 mm from the photon beam axis in the direction of expected plasma acceleration.

Figure 3-16. Electron plasma density measured by a Langmuir probe located immediately below the photon beam (i. e., in the vapor flow direction).

Further evidence of this enhanced directionality is shown in Figures 3-17 and 3-18, where oscilloscope traces are presented of the time-resolved ion current reaching the Faraday cup (positioned as shown in Figure 2-11). Figure 3-17 is a double exposure; the lower-amplitude trace represents ion current collected when the coil is not fired (i.e., the current due to free ambipolar expansion of the plasma column). In the higher-amplitude trace, the coil has been fired, and the total charge collected, $Q = \int i dt$, exhibits a substantial increase: $Q_{\text{coil}}/Q_{\text{ambipolar}} = 1.7$. A later mean arrival time of ions in this trace can be attributed to the fact that coil-firing is triggered by the arrival of the photon burst at a detector so that the greatest effect of the coil takes place late in time.

In Figure 3-18, the Faraday cup has been moved farther away from the plasma column ($z \approx 50\text{mm}$). The enhancement of charge collected when the coil is fired remains evident in these photographs, but is less dramatic ($Q_{\text{coil}}/Q_{\text{ambipolar}} \approx 1.3$). At this distance from the coil, the magnetic field is sufficiently weak that electrons are no longer adiabatic, and the type of expansion has become two dimensional.

The total charge collected by the Faraday cup is 26.5 pC with the coil not fired; this number is in reasonable agreement with the value $Q = 16.1 \text{ pC}$ ($\approx 10.8\%$ of 149 pC) which is predicted on the basis of initialized plasma density and the Faraday-cup collection geometry, as described in an earlier section. The charge collected with the coil operational was 45.1 pC, which represents 18.4% of the 245 pC total charge ($245 \text{ pC} - 25.6 \text{ pC}/0.0108$).

4. Effect of the Magnetic Field on the Discharge Plasma

The reflex discharge plasma source produces a flowing plasma which attains a density (at the photon beam axis) which is variable from zero to about 10^{17} m^{-3} ; the electron temperature is typically 1 - 2 eV, and the plasma potential is close to the potential of the anode (maintained at chamber ground).

In the absence of current in the drive coil, the Faraday cup detects only a very low level of ion current. This fact is indicative of the relatively high degree of localization of plasma leaving the discharge source.

Figure 3-17.

Ion current collected by the Faraday cup placed 10 mm from the photoplasma column (double exposure). The larger-amplitude trace shows ion current when the drive coil is fired, and the lower amplitude trace corresponds to free expansion of the plasma column (coil not fired).

Figure 3-18. Ion current collected by the Faraday cup. (a) Magnetic piston coil not fired. (b) Magnetic piston coil fired. Both traces were triggered on the leading edge of the photon burst.

When the drive coil is fired, a pulse of ion current is detected by the Faraday cup (typically with a peak current of 1 to 2 μA) followed by a long "tail" of low-level plasma which reaches the cup late in time (i. e., during the decay of the magnetic field). Figure 3-19 is a multiple exposure, showing current collected by the Faraday cup at three different distances from the plasma stream. The upper portion of this Figure shows the time-rate-of-change of the coil magnetic field, $\dot{B}$, as detected by a loop of wire placed near the coil. At $t = 16 \mu\text{s}$, $\dot{B}$ goes to zero, and the magnetic field has reached its maximum value. Analysis of the Faraday cup characteristics reveals that the leading portions of the current pulse propagate with a constant velocity of $\dot{z} = 1750 \text{ ms}^{-1}$. This velocity is consistent with one-dimensional ambipolar expansion of the plasma, assuming an electron temperature of $T_e = 2 \text{ eV}$. The action of the drive coil in this experiment is therefore explainable on the basis that transverse magnetic field lines leaving the coil isolate and capture a region of flowing plasma, causing it to diffuse with high directionality into the Faraday cup.

5. Effect of Magnetic Field Phase Angle on Charge Collection

Data presented in the previous sections demonstrate that the expansion and flow properties of plasmas are altered by the action of the drive coil magnetic field. It remains to be shown, however, whether the altered directionality results from active acceleration of the plasma by a $\underline{J} \times \underline{B}$ or $\mu \nabla B$ force, or from a change from two-dimensional to one-dimensional ambipolar diffusion. In the following discussion, data are presented which show the effect of varying the time-sequence of firing the photon source and drive coil; the photo-plasma can be created in a rapidly rising magnetic field, a slowly decaying one, or in one that is almost static.

A conceptual interpretation of the effect of delay time is shown in Figure 3-20, where multiple exposures of current detected by the Faraday cup have been made with the photon source fired at varying delay times after firing of the drive coil. The magnetic field, which is the same for each exposure in Figure 3-20, rises to its maximum

Figure 3-19.

Lower trace shows ion current collected by the Faraday cup at three positions: $z = 11, 18,$ and 24 mm . Upper trace shows time-rate-of-change of the drive coil magnetic field, $\dot{B}$, measured with a pickup loop placed near the coil.

Figure 3-20.

Effect of magnetic field phase angle on charge collected by the Faraday cup (multiple exposure). Upper trace: photon flux, each photon burst with a different delay time after drive coil firing. Lower trace: ion current collected by the Faraday cup, consisting of discharge plasma and photoplasma arrivals.

value at $t = 15 \mu s$, then slowly decays, reaching 10% of its peak value at $t = 25 \mu s$ (see Figure 3-19). The upper trace in Figure 3-20 shows the timing of photon emission. In the lower trace, two components of ion current are visible: a broad hump, extending from about $t = 3 \mu s$ to $t = 20 \mu s$ records the arrival of plasma from the DC discharge. In addition, spikes of ion current due to the photoplasma are visible; as the photon burst is delayed less and less, these spikes appear to slide to the left of the photograph, finally adding to the ion current from the discharge plasma. The dominant feature of these data is that the photo plasma spikes are of almost constant area (i.e., charge): only a slight enhancement is visible in the region of largest $\dot{B}$.

A more quantitative set of data characterizing this effect is presented in Figure 3-21, where the photoplasma charge Q collected by the Faraday cup is plotted against delay time. In these data, the charge due to the discharge plasma has been subtracted from the total, leaving only photoplasma charge. Although these data clearly indicate the substantial increase in charge collection brought about by the magnetic field, they do not support a claim that the charge enhancement depends on the time variation of the magnetic field, $\dot{B}$. The dominant importance of the magnitude of the field strength is illustrated in Figure 3-22. In this photograph, the photon source is fired at a time when $\dot{B}$ is negative (at a larger delay time than shown in Figure 3-21). The multiple exposures show that the charge collected increases and saturates as the drive coil (peak) current is increased from zero to 1.5 kA in five steps (the last three traces appear merged). If the $\underline{J} \times \underline{B}$ plasma acceleration mechanism were the dominant force here, increasing the coil current would tend to accelerate the plasma away from the Faraday cup, reducing the charge collected. This effect is opposite from the one observed, and it shows that the dominant parameter for charge collection is the magnitude of the field strength, B , rather than the time-rate of change of the field strength, $\dot{B}$.

Figure 3-21. Effect of magnetic field phase angle on photo-plasma charge, Q_Y , received by the Faraday cup.

Figure 3-22.

Effect of varying magnetic field strength on charge collected by the Faraday cup in a decaying magnetic field ($B < 0$).

Upper trace: photon flux. Lower trace: five exposures of the Faraday cup current taken with peak drive coil currents of 0, 500, 1000, 1250, and 1500 A. The amplitude of the traces increases with increasing coil current.

SECTION 4

DISCUSSION OF RESULTS

Data presented in the preceding sections support the conclusion that active acceleration of plasma by $\underline{J} \times \underline{B}$ or $\mu \Delta B$ forces has not been demonstrated. These data do conclusively show, however, that ambipolar expansion in a magnetic field is substantially more directional than free expansion. To the extent that the magnetic field geometry employed in these experiments is not intended to maximize plasma collection by ambipolar diffusion, the demonstrated collection efficiency $Q/Q_{\text{total}} = 0.18$ is remarkable.* Furthermore, the transport of charge to the Faraday cup is sufficiently rapid that charge-exchange losses might be expected to be minimal, even with $\tau_{\text{cx}} = 5 \mu\text{s}$. That ambipolar expansion works this well is evidently a result of good (near-classical) confinement of plasma electrons by the magnetic field. The saturation of collected charge with increasing magnetic field strength which has been demonstrated (cf Figure 3-22) may mean that very modest magnetic fields are adequate to produce one-dimensional

* The value $Q/Q_{\text{total}} = 0.18$ is not, however, an unreasonable value to expect. In a previous section, the charge collection was shown to be consistent with a purely geometrical model

$$\frac{Q}{Q_{\text{total}}} = \frac{1}{8} \frac{D \tan^{-1} \left(\frac{D}{2Z} \right)}{L + Z} = 0.108,$$

where D is the Faraday cup diameter, and L is the photoplasma column length. Consider now a model of ambipolar expansion in the drive coil field. Assume that expansion is uniform and two-dimensional in the yz plane. The geometry for this situation results in

$$\frac{Q}{Q_{\text{total}}} = \frac{1}{2\pi} \frac{D \tan^{-1} \left(\frac{D}{2Z} \right)}{L + Z} = 0.138.$$

To achieve higher Q/Q_{total} with ambipolar expansion in a magnetic field, most of the field lines which intersect the photoplasma must also intersect the Faraday cup.

expansion. These results are extremely encouraging from the standpoint of effective isotope separation, and beg further investigation. (A discussion of possible experimental approaches is given in Section 5.)

The absence in these experiments of active $\underline{J \times B}$ or $\mu \nabla B$ plasma acceleration may result from one or both of two causes: the values of B and $\dot{B}$ employed were simply too weak, and the passive extraction by one-dimensional ambipolar expansion was so rapid that this effect strongly dominated weaker active forces. The magnitude of the $\underline{J \times B}$ extraction force that was calculated earlier in this section for the drive coil employed for these experiments could be expected to result in a substantial acceleration of plasma. However, the plasma expands out of the extraction region at a speed that leaves little time for the $\underline{J \times B}$ force to act. It is not known what energy will be left in the photoelectrons in a real isotope-selective laser-ionization device. If the electrons are cold, ambipolar expansion will be slow, and $\underline{J \times B}$ more readily apparent; if they are hot, ambipolar expansion in a suitable magnetic field may be adequate in itself. In any event, both approaches merit further study and development, and a synergistic combination has special attractiveness.

In addition to the conclusions just reached regarding plasma extraction mechanisms, this program has produced several significant results for the simulation of plasma-separation processes:

- The dependence of the photoplasma density n_i^0 on neutral density has been established for both mercury and argon and shown to be consistent with expectations.
- The collection of freely expanding photoplasma over reasonable distances is consistent with purely geometric considerations.
- The expansion velocity of the photoplasma coincides with the ion-acoustic speed v_a at the measured electron temperature of $T_e = 1 \text{ eV}$.
- A strong directionality is imparted to both the discharge and photo plasmas by a very modest magnetic field.

The encouraging nature of the results summarized above suggests both the need for⁴ and the potential success of a continuing investigation of plasma extraction processes. Some specific recommendations for this investigation are presented in the following section.

SECTION 5

SUGGESTIONS FOR FUTURE WORK

This program has created an effective plasma-extraction simulation facility and data base from a number of diverse technologies. These technologies have been assembled in such a manner as to produce repeatable, self-consistent, and comprehensive data. Experiments conducted to date, however, have established only a rudimentary level of plasma-extraction technology. This technology can be advanced by several specific experimental approaches:

- Limits of ambipolar diffusion. The positive results achieved in this program should be followed up by investigating a magnetic field geometry in which all the field lines which pass through the photoplasma also pass through a "product collector" Faraday cup. Gating of the magnetic field to achieve high collection efficiency for a specific time interval should be pursued in this context.
- $J \times B$ revisited. A revised $J \times B$ extractor should be investigated; it should feature separate B_x and B_z field sources (notice that in the force equations, $J \times B$ involves only B_x and B_z). A bidirectional B_x (i. e., $B_x > 0$ for $x > 0$ and $B_x < 0$ for $x < 0$) would be established with dc field coils or permanent magnets. An improved magnet pulse circuit should also be developed, to provide larger B and $\dot{B}$ (up to the ionization limit discussed previously).
- Improved simulation. Means should be investigated to provide the flowing background plasma with a drift velocity independent of the thermal velocity (possibly by extraction and neutralization of an ion beam). This improvement would permit more realistic experimentation with magnetic gating, upstream deflection of photoplasma, and similar areas.
- Electron heating. The plasma extraction method that utilizes one-dimensional ambipolar expansion in a magnetic field is sufficiently rapid for effective isotope separation only if the electron temperature is high enough. Methods for supplementary electron heating (e. g., RF heating with microwaves or electron-cyclotron resonance heating) should be studied to prepare for the eventuality that the laser selective-ionization process leaves too little energy in the photoelectrons.

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