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High-Power Pulsed-Ion-Beam Acceleration and Transport

Stanley Humphries, Jr, Thomas R. Lockner

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TABLE OF CONTENTS

	<u>Page</u>
I. INTRODUCTION	1
II. CHARACTERISTICS OF ION BEAM NEUTRALIZATION	3
A. Space Charge Limitations on Ion Beam Acceleration and Transport	3
1. Acceleration Gaps	3
2. Periodic Lens Arrays	11
3. Propagation in Free Space	15
B. Methods of Neutralization	16
III. TRANSPORT OF INTENSE ION BEAMS	23
A. Transverse Neutralization	23
B. Longitudinal Neutralization	33
C. Focusing Intense, Neutralized Ion Beams	41
D. Ion Beam Transport by Electron Control	56
IV. HIGH CURRENT ION BEAM ACCELERATOR THEORY	64
A. Magnetically Insulated Acceleration Gaps	64
B. Multi-Gap Effects in High Current Ion Accelerators .	79
1. Transverse Beam Control	80
2. Longitudinal Beam Control	88
V. EXPERIMENTS ON INTENSE ION BEAM ACCELERATION	97
A. Ion Sources	97
B. Electron Sources	101
C. Ion Injectors	103
D. Post-Acceleration and Transport	104
E. Inertial Fusion Accelerators	110
VI. CONCLUSIONS	113
REFERENCES	117

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Stanley Humphries, Jr. and Thomas R. Lockner
Sandia National Laboratories
Albuquerque, NM 87185

I. INTRODUCTION

In the past eight years, methods have become available for the generation and transport of very high current pulsed ion beams (see, for instance, Humphries, 1980a). These methods are based on a new understanding of electron neutralization in vacuum regions. The addition and control of electrons can greatly reduce space charge effects in ion beams on fast time scales. Almost all work in this field has been directed towards the construction of extremely high current, single-stage, single-shot devices (ion diodes) (Dreike, et al., 1979; Goldstein, et al., 1978; Miller, et al., 1979). In contrast, efforts in the Pulselac Program (Humphries, 1978a) at Sandia National Laboratories have concerned understanding the basic processes and limitations of neutralization, developing practical injector technology, and utilizing neutralized transport for multi-stage acceleration. At present, this is the sole program directed toward multi-stage systems for intense ion beams.

Primary interest in high intensity pulsed ion accelerators has derived from applications to inertial confinement fusion. The technology developed has advanced the reliability of intense ion injectors to a point where they can be considered for other

applications, such as pulsed neutron radiography, heating magnetic confinement fusion experiments, and machining and annealing microcircuits. In this paper, the state of research on intense ion beam acceleration and transport will be reviewed. Section II summarizes the limitations imposed on ion beam transport by space charge effects and methods available for neutralization. The general problem of ion beam neutralization in regions free of applied electric fields is treated in Section III. The physics of acceleration gaps is covered in Section IV. Finally, experiments on multi-stage ion acceleration are summarized briefly in Section V. Expressions throughout the paper are given for singly charged ions in CGS units.

II. CHARACTERISTICS OF ION BEAM NEUTRALIZATION

Space charge repulsion is the main impediment to the generation and propagation of high current ion beams. A variety of methods have been developed to reduce the effect of space charge. Methods of neutralization will be reviewed in this section. To begin, the effects of space charge on bare beams will be given. This will define the magnitude of the problem and will serve to introduce material of use in latter sections.

A. Space Charge Limitations of Ion Beam Acceleration and Transport

Space charge repulsion is a particular problem for bare ion beams, which move slower than electron beams of the same energy. Longitudinal space charge limits occur when electrical potentials are comparable to the beam energy. Transverse limits result when space charge forces exceed the ability of a focusing system to contain the beam. Longitudinal space charge effects are important in the first gaps of a multi-stage system; transverse limits are of concern throughout the acceleration and transport processes.

1. Acceleration Gaps

Ion beams are generally unneutralized in acceleration gaps, even if methods of neutralization can be used in other parts of the accelerator. Longitudinal limits set by space charge will be considered first. In the one-dimensional case illustrated in Figure 1, ions with an initial energy $[(m_i v_i^2/2) - eV_0]$ enter a gap

a)

b)

Figure 1. Longitudinal space charge limits for one-dimensional ion flow. (a) Geometry and parameters for calculation. (b) Allowed ion flow normalized to the Child limit, j_{ic} . The parameter β is the ratio of the ion energy exiting the gap to eV_0 .

with an applied voltage of V_0 . In the absence of space charge the gap potential varies linearly from eV_0 to 0; with self-field effects, higher potentials can occur. In particular, the maximum potential allowed is $m_i v_i^2/2$; this defines the limiting flux of ions. Poisson's equation can be solved with the boundary conditions $\phi(0) = V_0$, $\phi(d) = 0$, $d\phi/dx = 0$ when $e\phi = m_i v_i^2/2$, and an ion density of the form $n_i(x) = n_i(d)(1 - 2e\phi/m_i v_i^2)^{-1/2}$. The limiting current density can be expressed in the form

$$j_i \leq j_{ic} F(\beta) , \quad (1)$$

where

$$j_{ic} = \frac{1}{9\pi} \sqrt{\frac{2e}{m_i}} \frac{V_0^{3/2}}{d^2} .$$

The function $F(\beta)$, the enhancement factor for post-acceleration gaps, is given by

$$F(\beta) = \beta^{3/2} [(1 - \beta^{-1})^{3/4} + 1]^2$$

where $\beta = m_i v_i^2/2eV_0$. The quantity j_{ic} is the familiar Child limit (Child, 1911) for a gap with zero initial ion velocity. The function $F(\beta)$ is plotted in Figure 1(b). Longitudinal space charge effects are important only in the first gap, which we will refer to as the injector gap. In a constant current density system, space charge potentials have a rapidly decreasing relative effect as the beam is accelerated. Methods for alleviating space charge limitations in injector gaps will be discussed in Section IV,A.

Space charge fields cause transverse defocusing of beams unless compensating forces are applied. The fields associated with space charge in an acceleration gap can be calculated analytically for the case shown in Figure 2(a). The gap between two conductors has spacing d , the beam has width w , and the ion density, n_{i0} , is uniform. This is a good approximation for a post-acceleration gap where the relative change in ion velocity is small, and hence n_i is almost uniform in the x -direction. The space charge potential is

$$\frac{\phi(x,y)}{\phi_0} = \frac{16}{\pi^3} \sum_{m=1,3,5,\dots}^{\infty} \frac{1}{m^3} \sin\left(\frac{m\pi y}{d}\right) x \left\{ 2 - e^{-m\pi x/d} - e^{-m\pi(x-w)/d} \right\} \quad (2)$$

where $\phi_0 = \left(\frac{\pi}{2}\right) en_{i0} d^2$.

The scaling potential, ϕ_0 , is the maximum for a one-dimensional beam ($w/d \gg 1$). A plot of this quantity is given in Figure 2(b). The quantity $E_x(d/2, x)$ is graphed in Figure 2(c). Finally, Figure 2(d) shows typical beam density as a function of current density and energy. Ion beams with current density in the characteristic range 10-100 A/cm² have high density, exceeding 10¹² cm⁻³ at low energy. Transverse field effects in acceleration gaps can be significant.

Figure 2. Effects of space charge in an acceleration gap bounded by two conducting planes. (a) Geometry and parameters, showing uniform ion density of finite extent in x . (b) Normalizing potential, $\phi_0 = \pi n_{i0} d^2 / 2$. (c) Variation of transverse electric field at the center of the gap as a function of w/d . (d) Typical values of proton beam density as a function of beam energy and current density.

Figure 2(b).

c)

Figure 2(c).

d)

Figure 2(d).

The plot of Figure 2(c) illustrates an important method of reducing space charge effects, neutralization by image charges in conductors. For low w/d , the transverse electric fields increase almost linearly away from the beam axis. At higher w/d , image charges in the electrodes are important, and the variation is non-linear. For wide beams, space charge fields occur within a distance d from the edge; the transverse field reaches a maximum value of $1.5 (\phi_0/d)$. Thus, by increasing the width of beams passing through conducting electrodes, higher total currents can be accelerated without an increase in transverse defocusing. Of course, in practice, physical conductors cannot be located in the beam path. The usual solution is to divide the beam into many small beams that pass through holes in conducting electrodes. In such multi-aperture systems, the holes must have diameter $w \leq d$ to limit defocusing forces; the presence of image charges on the conductors isolates each small beam. The space charge forces acting on each small beam are reduced from the cumulative effect of the full beam.

2. Periodic Lens Arrays

Ion beams are usually contained in accelerators by a series of focusing lenses. Confinement against space charge forces may be performed by a separate series of conventional charged particle lenses or by the addition of transverse field components to the accelerating field in each gap (see, for instance, Brewer, 1967). In this section, a model for dealing with space charge effects in periodic lens arrays by means of average

confinement energy potentials will be developed. This model will be used in Section IV,B,1 to treat particle confinement in highly non-linear focusing systems.

An array of thin lenses is shown in Figure 3. For a linear lens, it is possible to define a focal length, f . For lenses with uniform spacing ℓ , particle orbits will be stable if $f > \ell/2$ (see, for instance, Ramo, et al., 1965). Although there is a straightforward solution to the orbit difference equations, we will consider the limit in which the axial distance travelled during a transverse oscillation is much greater than ℓ . In this limit, focusing forces act almost continuously, and can be expressed in terms of an average focusing force. This force can be spatially integrated to yield an average potential energy. Particle orbits are oscillations in this potential energy well. Various focusing and defocusing potentials can be added and self-consistent, stationary phase distributions can be easily derived (Humphries and Poukey, 1979). It will be shown in Section IV,B,1 that the condition for the validity of this model also guarantees that beams will propagate with small emittance growth, even with non-linear lenses.

For a thin lens, the average focusing force is $F_f = m_i \Delta v / \Delta t$, where Δv is the change in transverse velocity in passing through the lens and $\Delta t = \ell / v_i$. Some forces in the accelerator (such as space charge in acceleration gaps) may provide a known force that acts periodically over a certain distance, d . In this case, the average force is simply the known force multiplied by the fraction of time it acts, (d/ℓ) .

Figure 3. Periodic series of linear lenses for beam confinement.

For the case of Figure 3, the approximate equation of motion is

$$\frac{d^2 r}{dz^2} = -\frac{r}{f\ell} \quad (3)$$

with solution

$$r = r_0 \cos \left[\frac{z}{\sqrt{f\ell}} + \delta \right]$$

The average focusing potential energy is quadratic, given by

$$U_f(r) = \left(\frac{m_i v_1^2}{2} \right) \left(\frac{r^2}{f\ell} \right)$$

The average potential energy for space charge forces acting continuously on a uniform density cylindrical beam is given by

$$U_e(r) = -\pi e^2 n_{i0} r^2$$

In order that a beam with a maximum transverse velocity, v_{rmax} , be contained by the lens system, the condition

$$U_f(r_{max}) \geq U_e(r_{max}) + m_i v_{rmax}^2 / 2 \quad (4)$$

must be satisfied. This defines the space charge limit for transport. As an example, consider motion of proton beams through a series of solenoidal magnetic lenses filling a fraction ϵ of the accelerator. The effective potential energy for such a lens is

$$U_f(r) = \epsilon \left[\frac{e^2 B_0^2}{8 m_i c^2} \right] r^2$$

where B_0 is the magnitude of the solenoidal field. The space

charge limit for a zero emittance beam is

$$\left(\frac{B_o^2}{8\pi}\right) \geq n_{i0} m_i c^2/\epsilon$$

If $B_o = 40$ kG, the maximum beam density is $4.2 \times 10^{10} \text{ cm}^{-3}$. Consulting Figure 2(d), this density corresponds to low values of current density. Although the situation is improved in strong focusing lenses, it is generally true that magnetic focusing is ineffective for low velocity ions. Calculations on transport in strong focusing lens arrays using the continuous approximation have been carried out (Courant and Maschke, 1976; Reiser, 1977; Reiser, 1978).

3. Propagation in Free Space

Many applications call for propagation of an ion beam through a vacuum region. Longitudinal limits can be significant for beams produced in a space charge limited single stage injector. For instance, if a large cross section ion beam is produced in an injector with gap spacing d , then the ions cannot propagate in vacuum more than a distance $2d$. For high energy beams, longitudinal limits are generally not a problem since the maximum space charge potentials are less than the beam energy. In this case, transverse space charge effects are of greatest concern.

Transverse spreading of uniform density beams can be treated easily in the laminar approximation (see, for instance, Wilson and Brewer, 1973). Since electric fields grow linearly away from the beam axis, it is sufficient to follow the equation

of motion of the beam envelope. For a cylindrical beam with orbits parallel to the axis, velocity v_i , initial density n_{i0} and initial radius r_0 , the paraxial envelope equation is

$$\frac{d^2 r}{dz^2} = \frac{r}{2} \left[\frac{4\pi e^2 n_{i0} r_0^2}{m_i v_i^2} \right]. \quad (5)$$

This yields

$$\int_1^{r/r_0} \frac{dR'}{\sqrt{\ln R'}} = \sqrt{\frac{4\pi e^2 n_{i0}}{m_i v_i^2}} z \quad (6)$$

The solution to Equation (5) for initially parallel beams is plotted in Figure 4. The expansion of a slit beam of initial width w_0 , given by $w(z)/w_0 = 1 + (4 e^2 n_{i0}/m_i v_i^2)z^2$ is also plotted. As an example, 1 cm diameter, 10 A beam of 200 keV protons will double its width in only 2.7 cm. The conclusion is that, except at very high energy, ion beams cannot be propagated a significant distance in vacuum without neutralization.

B. Methods of Neutralization

A variety of methods have been developed to add negative charge to ion beams to reduce space charge effects. A categorization of these methods is given in Figure 5. One approach that has already been mentioned is neutralization by image charges. In this case, electrons are bound within and travel through conductors. This process underlies multi-aperture

Figure 4. Expansion of initially parallel, unneutralized ion beams in vacuum. The quantity n_{i0} is the initial ion density, r_0 is the initial radius (for a cylindrical beam), and w_0 is the initial width (for a slab beam).

NEUTRALIZATION METHODS

Figure 5. (No caption).

injectors developed for space propulsion (see, for instance, Stuhlinger, 1964) and neutral beam production (see, for instance, Kunkel, 1979). Multi-aperture, multi-stage systems have recently received attention for inertial confinement fusion driver applications (Herrmannsfeldt, 1978; Adams, et al., 1980; Guiragossian, 1980). Accelerators with image charge neutralization operate with bare beams and localized electrons; manufacturing difficulties limit the homogeneity of positive and negative space charge that can be achieved. Therefore, these accelerators are limited to medium current density ($< 1 \text{ A/cm}^2$). For high current operation, free electrons must enter the ion beam volume.

Free electron neutralization can be divided into plasma and vacuum categories, depending on whether electrons are initially present in the transport region. In plasma neutralization, the beam passes through a background of ions and electrons with density typically much greater than that of the beam. The plasma in the transport region can either be preformed or formed from the collisions of the beam with a background gas. The latter case has important application to intense neutral beam injectors (Holmes, 1979); the low energy beams are neutralized after emerging from acceleration gaps by a plasma formed in the residual gas from the neutralization cell. When the background gas is ionized, the ion escapes while the electron remains within the ion beam.

The response time of this neutralization process is limited by either the collision rate or the escape time of the low energy ion (Mazarakis, et al., 1979). It is thus useful only for long pulse ($\gg 1 \mu s$) or continuous beams.

Neutralization in a preformed plasma can be rapid if the plasma is electrically connected to a source of electrons. In this case, electrons are added to the beam volume rather than removing background ions. Plasma electrons respond to a change in beam space charge in a time scale $1/\omega_{pe}$, where ω_{pe} is the plasma frequency, $\omega_{pe} = \sqrt{4\pi e^2 n_p / m_e}$. This is less than 1 ns for typical plasma density $10^{12} - 10^{14} \text{ cm}^{-3}$. Since plasmas are good conductors, they can also bring about net current cancellation of pulsed ion beams by inductively driven return currents. Over times less than the current decay time of the plasma column occupied by the beam, the electron return current has the same spatial distribution as that of the beam current. At longer times, resistive processes spread the electron current over a larger cross-section. Thus, self-magnetic field effects can occur even if the net current remains zero. Plasma transport of pulsed ion beams has been studied extensively for application to heavy ion transport in inertial fusion reactors (see, for instance, Olson, 1980). An interesting option, studied for light ion fusion applications, is to imbed an applied magnetic field in a plasma before entrance of the beam (Miller, et al., 1977; Goldstein, et al., 1977). This field can couple the transverse momentum of the beam to the inertia of the surrounding plasma. In this

case, the beam can be simultaneously space charge neutralized, current neutralized, and confined for transport.

Although plasma neutralization can be effective, there are a number of reasons why it would be impractical to utilize the process in high current accelerators. First, it is difficult to produce a sufficiently dense plasma over a large transport volume. Second, the plasma must be completely isolated from the acceleration gaps to avoid gap shorting or acceleration of plasma particles. Nonetheless, it must be close enough to the gap to provide effective neutralization. Finally, charge changing collisions of the beam with the plasma would limit the possible accelerator length. This could be less than 1 meter for heavier ions.

These problems can be solved by using vacuum neutralization. In this process, both beam ions and neutralizing electrons are introduced simultaneously into the propagation volume. The easiest way to accomplish this is to allow electrons to be drawn in by the space charge fields of the ion beam. This process (which we will refer to as auto-neutralization) proceeds until the beam fields are cancelled, resulting in a close balance of positive and negative space charge. Utilizing auto-neutralization and understanding its limitations with respect to time response and other factors is the main subject of this paper.

Vacuum neutralization can be subdivided according to whether an applied electric field is present or not (acceleration versus transport). These categories can be further

divided into longitudinal and transverse neutralization, depending on whether electrons enter colinearly with the ions or are introduced from the periphery of the beam. In acceleration gaps, the longitudinal and transverse cases represent the two major methods of electron control in the presence of applied electric fields, reflexing and confinement by transverse magnetic fields (magnetic insulation). The latter case, the magnetically insulated gap, will be discussed in detail in Section IV,A. The reflex triode (Creedon, et al., 1975; Prono, et al., 1975) provided the first demonstration that ion beam flow in acceleration gaps could become independent of conventional space charge limits. The problem of transport is discussed in the next section. This area has been studied previously for applications to space propulsion. In this early work, neutralization was treated by analytic one-dimensional equilibrium models. These models implied that auto-neutralization would result in a poor mixing of ions and electrons. The main thrust of the recent work discussed in this paper has been to investigate the time-dependence of auto-neutralization and to demonstrate that in real systems without one-dimensional symmetry homogeneous mixtures of externally produced ions and electrons can be formed.

III. TRANSPORT OF INTENSE ION BEAMS

Theories of vacuum auto-neutralization will be considered in this section. Much of this work has evolved from observations of 10-100 kA beams produced by high current ion diodes. Transverse neutralization is treated first. This has application to the transport of ions between acceleration gaps of high current accelerators and to the passage of intense beams through conventional magnetic lenses. The second case is longitudinal neutralization, where electrons enter colinearly with the ions when the beams pass through foils, grids or plasma clouds. The major application of this process is in vacuum transport of beams after leaving the accelerator. In the third part, lenses for intense ion beams are treated. The section concludes with a general discussion of electron control for high current ion beam transport.

A. Transverse Neutralization

A typical geometry for transverse neutralization is shown in Figure 6(a). While ions move longitudinally, electrons can move only in the transverse direction. This may come about from the symmetry of the problem or through the application of transverse magnetic fields. If the electron sources can supply a space charge limited flow, the condition $E_x = 0$ holds on the boundaries. This condition implies that the total number of ions and electrons is equal; the main problem is in determining whether they are distributed homogeneously. In most neutralization problems, electrical potentials are of the order of

Figure 6. Transverse neutralization. (a) Geometry. (b) Analytic equilibrium solutions for a pure one-dimensional case with ions present for all time. Dashed line: Potential with no electrons, uniform ion density. Solid line: Lowest order solution with space charge limited electrons.

$(m_e/m_i)E_i$, where E_i is the energy of the ions. While these potentials can result in large variations in electron density, the ion density can be considered unaffected by electric fields.

For a uniform ion density and an electron density given by $n_e(x) = A\phi^{-1/2}$, Poisson's equation can be solved (Stuhlinger, 1964) in the geometry of Figure 6(a) with the boundary conditions $E_x(-d/2) = E_x(d/2) = 0$. There is an infinite sequence of solutions, characterized by mode number m . The potential midway between boundaries is given by

$$\phi(x=0)/\phi_0 = \frac{4}{[(2m+1)]^2} \quad (7)$$

where ϕ_0 was defined in Equation (2). The solution with $m=0$ is considered the most likely to occur in practice. Solutions with and without electrons are plotted in Figure 6(b). The central potential is reduced only a factor of 2.5 with electrons. This comes about because the electron density diverges at the boundaries, and the beam is under-neutralized at the center. Two modifications must be made in order to calculate what happens in a real physical situation (Humphries, et al., 1976; Humphries, et al., 1977; Humphries, 1978b). First, the equilibrium model resulting in Equation (7) has the inherent assumption that the electron and ion densities existed for all previous time. Processes occurring during the build-up of ion density must be included. Second, no real system will have perfect one-dimensional symmetry. With two-dimensional effects,

b

Figure 6(b).

electrons can develop velocity components in the y and z directions which will significantly alter the density divergence at the boundaries.

Computer simulations have been performed with a time-varying ion density and two-dimensional electron motion (Poukey and Humphries, 1978). The width of the transport region was 1 cm. The spatially uniform ion density rose to its final value 10^{12} cm^{-3} in 0.5 ns. A magnetic field was included which could be orientated along x (for the ideal one-dimensional case) or tilted at an angle to introduce additional velocity components to the electron motion. In real systems, these effects could arise from curved magnetic fields, irregular boundaries, or longitudinal variation of ion density. Again, electron emission was space charge limited. Typical results for the variation of central potential with time are shown in Figure 7. Case A is for purely transverse magnetic field. The potential rose to a steady state value, somewhat reduced from the prediction of Equation (7) because of the removal of electrons from the boundaries by the rising beam space charge fields. This solution is still inadequate for neutralized transport. The situation changes markedly for Case B in which the magnetic field was tilted 15 degrees with respect to the x-axis. With the non-collisional, elastic randomization of the electron velocity distribution, the potential collapsed after the rise of ion density, indicating almost complete neutralization. Similar effects are evident

Figure 7. Transverse neutralization solutions with time-dependent introduction of ions and two-dimensional effects. Results of computer simulation show central potential versus time. Top dashed line: Potential from lowest mode of simple analytic solution. Dashed line: Time variation of ion density. Curve A: Purely one-dimensional solution with ion temporal variation shown. Curve B: Solution with small two-dimensional effects added. Also shown is the upper limit on potential during the rise of ion density set by electron flow, $\phi_0 \text{ max}$.

in actual two-dimensional simulations (see Section IV,A). The calculation illustrates an important aspect of treating vacuum neutralization theoretically. In the initial state, electrons and ions are separated; in the final state they should be mixed homogeneously. The neutralization process must be irreversible, and theoretical treatments must allow pathways for relaxation of the electron distributions to demonstrate good mixing.

Limitations on the time response of transverse neutralization are important for application to pulsed ion beams. The two main impediments are the flow rate of electrons into the beam volume and the relaxation rate of the electron distribution. In most pulsed beam experiments, strong longitudinal electric fields can occur if the ion beam is unneutralized, so that electron randomization can take place rapidly. This may be aided by the electron-ion two-stream instability (McBride, et al., 1972), which thermalizes the electron distribution. The limit imposed by space charge flow of electrons into the beam has been estimated (Humphries, 1978b) by finding the central potential needed to balance the rate of change of electron and ion space charge when the ion beam rises to a value n_{i0} in a time Δt_r . For the geometry of Figure 6(a), the central potential is given by

$$\phi \lesssim d^2 \left(\frac{9\pi}{8} \frac{e n_{i0}}{\Delta t_r} \sqrt{\frac{m_e}{2e}} \right)^{2/3} \quad (8)$$

The upper limit occurs if all the ion charge were concentrated at the center of the transport region. If the ion charge is distributed closer to the source, the required potential is lower. The prediction of Equation (8) is plotted in Figure 7. The actual potential observed during the rise of ion density in Case B is reduced about a factor of 3 because the ion density was distributed uniformly.

We conclude this topic by considering some of the problems that may be associated with transverse neutralization. Equation (8) implies that large potentials can occur for vacuum transport of very high current pulsed beams from ion diodes (> 100 kA). Inclusion of a plasma with density comparable to or greater than that of the beam can greatly alleviate this problem. In this case, the d^2 factor in Equation (8) is replaced by λ_d^2 , where λ_d is the plasma Debye length. For typical plasmas, λ_d can be much smaller than 1 cm (see Section III,D). Another aspect of the transverse neutralization process is that of current cancellation. If a transverse magnetic field is present, extremely high electric fields are necessary for direct longitudinal electron motion ($E_z > B_x$). In most cases, electrons are stationary. Thus, the effects of the self-magnetic fields of the ion beam must be considered near magnetically insulated gaps and in magnetic lenses.

The stability of the ion-electron mixture in the case of transverse neutralization has begun to receive attention. Analyses indicate that an electron-ion two-stream instability can occur (Mondelli, 1981). This is quenched when the

electrons are thermalized to a temperature of $kT_e \lesssim (m_e/m_i)E_i$. Stability theory for the ions will be greatly modified from that performed for conventional accelerators. By cancelling space charge fields, neutralization removes a large source of free energy available to drive instabilities. In the transverse neutralization case, electrons make a direct electrical connection of the beam to the surrounding drift tube walls so that electric fields generated by instabilities can be shorted.

A final area of concern is whether collisions with the background neutralizing electrons can affect the ion charge state, bringing about beam losses. An advantage of vacuum neutralization for this process is that the electron density need be only as large as the ion density. As an estimate of the magnitude of this effect, we consider charge changing collisions of an N^+ beam accelerated to 46 MeV in an accelerator with an average longitudinal gradient of $d(eV)/dz = 1.5$ MV/m, typical parameters for inertial fusion applications. The cross section for the reaction ($N^+ + e^- \longrightarrow N^{++} + 2e^-$) is shown in Figure 8 (Harrison, et al., 1963). The probability of a reaction is $P = \int n_e \sigma dz$, where n_e is the background electron density. If the neutralized beam pulselength and area are constant, then the electron density varies approximately as $n_e = n_{e0} v_{i0}/v_i$, where v_i is the longitudinal ion velocity. For a constant average voltage gradient, $dz = m_i v_i dv_i / [d(eV)/dz]$. The reaction probability can be written

Figure 8. Typical cross section for charge changing collisions of ions with a background electron density.

$$P(v_i) = \frac{m_i n_{eo} v_{io}}{e(dV/dz)} \int_{v_{io}}^{v_i} \sigma(v_i') dv_i' \quad (9)$$

Performing the integral over the cross section, it is found that the probability for a reaction during acceleration to 46 MeV in a length of 31 meters is .018. In general, collisions with background electrons have little consequence for light ions, but may provide limitations to the acceleration of heavy ions to very high energy.

B. Longitudinal Neutralization

In the longitudinal neutralization process, electrons enter the propagation volume parallel with the ions, as shown in Figure 9(a). Electrons may be produced on a grid or foil or in a plasma cloud. Using multiple grids, it is possible to accelerate the electrons to a velocity equal to that of the ions. Treatments in this paper will be limited to the auto-neutralization case, where low energy electrons are accelerated by longitudinal space charge fields. An analytic equilibrium solution (Stuhlinger, 1964) can be found for the one-dimensional auto-neutralization process in the case where a uniform ion density (n_{i0}) extends to infinity from a source at $x = 0$. At the origin, $\phi(0) = 0$ and space charge limited electron flow implies that $d\phi/dx(0) = 0$. In equilibrium, the electron flux is independent of position, and hence $n_e = A/v_e = A(2e\phi/m_e)^{-1/2}$. With these conditions, Poisson's

Figure 9. Longitudinal auto-neutralization. (a) Geometry and boundary conditions for one-dimensional calculations. (b) Analytic solutions. Solid line: Steady state solution for ions and electrons filling the entire space, $x > 0$. Dashed line: Solution with electrons emitted while ions enter the space $x > 0$.

equation can be written in dimensionless form as

$$\frac{d^2\Phi}{dx^2} = \Phi^{-1/2} - 1 \quad (10)$$

where

$$\Phi = e\phi / \left(\frac{1}{2} m_e v_i^2 \right)$$

and

$$X = x / \sqrt{\frac{(m_e v_i^2)/2}{4\pi e^2 n_{i0}}}$$

The solution of Equation (10) is

$$X = \sqrt{2} \left\{ \sin^{-1}(\Phi^{1/2} - 1) - [\Phi^{1/2}(1 - \Phi^{1/2}/2)]^{1/2} + \pi/2 \right\} \quad (11)$$

Equation (11) is plotted in Figure 9(b). The potential oscillates between 0 and $4(m_e v_i^2/2)$ with spatial periodicity $(\pi m_e v_i^2/e^2 n_{i0})^{1/2}$.

This solution, with its oscillating downstream potential, again illustrates the difficulty in accounting for the irreversible process of neutralization with simple, symmetric models. As with the transverse problem, realistic effects must be added to demonstrate homogeneous mixing. For instance, the calculation of electron density assuming an existing uniform ion density neglects processes that could occur when the ions were filling the plane $x > 0$. An alternate analytic sheath solution has been derived for electrons extracted into a moving ion distribution, with the possibility of the reflection of electrons from a virtual cathode at the moving front

(Humphries, 1978c). The solution is identical to Equation (11) between the origin and the first potential maximum, but downstream the potential is constant, $\phi(x) = 4(m_e v_i^2/2)$, as shown in Figure 9(b). The electron velocity distribution is double-humped, with half the electrons entering at $v_e = 2v_i$ and half reflected from the front at $v_e = 0$. This distribution is unstable to the two-stream instability, which would thermalize the electrons.

As with the case of transverse neutralization, computer simulations provide a means to investigate the full problem with its time dependence and possible microinstabilities. Recent calculations (Humphries, et al., 1981b) have considered self-consistent electron motion in the presence of the space charge fields of a uniform density ion beam moving into the half plane $x > 0$ at velocity v_i . Electron phase space plots are shown in Figure 10(a) and (b). Near the source, the potential reaches $4(m_e v_i^2/2)$ across a sheath of about the same width as the analytic solutions. Downstream, the electron distribution is thermalized by the two-stream instability, as shown in Figure 10(c). The electron distribution is centered about $v_e = v_i$ with a half width (at half maximum) of $0.3 v_i$. The thermalization process causes the potential oscillations to damp out within a few sheath widths from the source, so that the downstream potential is approximately constant, $\phi(x) \cong m_e v_i^2/2$. The main conclusion is that complete neutralization is rapidly attained, and the background electron distribution will have a thermal spread about the beam velocity. In the case shown

Figure 10. Computer simulations of longitudinal auto-neutralization. (a) Electron longitudinal phase space plot when ions have moved a few sheath widths away from the source. Note oscillating potential, in agreement with Figure 9(b). (b) Electron phase space plot at a later time. Note broadening of initial sheath and thermalized downstream distribution. (c) Electron velocity distribution taken over region shown by dark line in part (b).

(with only one boundary), the solution also indicates net current cancellation. The one-dimensional auto-neutralization solution is applicable when the beam has transverse dimension large compared to the sheath thickness; this is equivalent to the requirement that the beam have high enough density to produce space charge potential in excess of $4(m_e v_i^2/2)$.

A charge and current neutralized ion beam in free space is stable. This can be easily demonstrated by considering a transformation to the rest frame of the beam. Ions are stationary and electrons have a Maxwellian distribution, so there is no source of free energy. There has recently been discussion (Poukey, et al., 1981b; Poukey, et al., 1981c) on whether neutralized beams in regions surrounded by conducting boundaries are unstable to the Pierce instability (Pierce, 1944). In the one-dimensional case of the instability, electrons are actively injected with velocity $v_e = v_i$ from one boundary through a uniform ion density. In a few electron transit times, analytic theory and computer simulations show a loss of neutralization, with a solution similar to that of Equation (7). The mechanism of the instability is as follows. If there is a slight underinjection of electrons ($n_e(0) < n_{i0}$), there will be a positive potential between the boundaries, which accelerates the electrons and lowers the density at the center. This process grows until the central potential is a substantial fraction of ϕ_0 . The inverse process can occur when there is an overinjection of electrons. The situation with two boundaries

and active electron injection seldom occurs in practice. We used computer simulations to investigate stability properties of beams in more common experimental situations. For active electron injection, we found that the beam must be electrically connected to both boundaries before the total number of electrons in the system can change substantially. There is thus no instability during the time it takes the ion beam to travel across the bounded region. For auto-neutralization, the $E = 0$ boundary conditions act as a feedback mechanism to maintain the total number of ions and electrons in the system equal. Thus, auto-neutralized beams propagating between boundaries are stable at all times.

Boundaries can affect current cancellation in auto-neutralized beams. In geometries, such as those of Figure 11(a), the electron return-current can decay over times longer than an ion transit time. A change in potential is necessary along the path shown to accelerate electrons to velocity $v_e \leq v_i$. In a conducting chamber, this voltage must be balanced by a change in magnetic flux. If the wall radius is more than twice the beam radius, the inductance acting on a net beam current is contributed mainly by the area outside the beam. In this case, a simple model of beam current growth can be derived by averaging quantities across the beam radius. If the beam is space charge neutralized, then $n_e \simeq n_i$. In order to accelerate electrons to velocity v_e , a voltage difference equal to $m_e v_e^2 / 2e$ must appear along the beam column. The net beam current is $I = I_b(1 - v_e/v_i)$, where I_b is the ion current.

Figure 11. Generation of net currents in auto-neutralized ion beams moving through regions bounded by conductors. (a) Geometry and parameters of calculation. (b) Normalized growth of net current ($I = I_e + I_b$) with time, where I_e is the neutralizing electron current, I_b is the ion current, and τ is defined in the text. Solution valid when $r_b < r_w/2$.

The voltage produced by changing external flux is

$V = - (1/c)(d\phi/dt) = (2l \ln(r_w/r_b)/c^2)(dI/dt)$. Equating the voltages gives an expression for the variation of net beam current with time

$$\frac{d}{d\tau} \left(\frac{I}{I_b} \right) = \left(1 - \frac{I}{I_b} \right)^2 \quad (12)$$

where $\tau = t/t_d$, and t_d is the characteristic decay time

$$t_d = \frac{4eI_b l \ln(r_w/r_b)}{m_e c^2 v_i^2}$$

The solution of Equation (12), $I/I_b = 1 - (1 + t/t_d)^{-1}$, is plotted in Figure 11(b). Decay times are long for high current beams in inductive chambers. For example, a 10 kA H^+ beam at 0.2 MeV in a chamber 1 meter in length with $r_w/r_b = 2$ has $t_d = 12 \mu s$. On the other hand, net currents can appear quickly for low current beams in small chambers; this process has potential use for high current beam diagnostics (Drobot, et al., 1978).

C. Focusing Intense, Neutralized Ion Beams

Applications such as inertial confinement fusion call for focusing neutralized ion beams to high power density. In some respects the problem is more tractable for neutralized beams than for space charge dominated beams since, to first order, the ions follow single-particle orbits in magnetic focusing elements. In this section, focusing of intense ion beams will

be treated, with emphasis on toroidal field lenses. After a consideration of focusing limitations set by lenses, the effects of neutralizing electrons on a focused ion beam during transit to the target will be summarized.

The transverse neutralization process described in Part A can occur in lenses that utilize transverse magnetic fields (Humphries, 1977b). Quadrupole lenses are ideal for neutralization since they have grounded pole pieces in close proximity to the beam. Electrons can be supplied from these surfaces. Solenoidal lenses are also a possibility, but have the disadvantages that focusing is second order and electrons must travel relatively long distances along field lines to enter the ion beam volume. Experiments by Maenchen (Maenchen, et al., 1979), in which a 60 ns ion beam passed through a solenoidal lens, have indicated space charge effects at the center of the lens consistent with Equation (8). Such space charge fields, brought about by the controlled introduction or starving of electrons, can be useful for ion beam transport and are discussed in Part D.

High current ion accelerators using magnetic field control of neutralizing electrons produce beams of annular cross section (see Section IV). For this reason, it is worthwhile to consider the toroidal field lens which can provide direct focusing of annular beams. This lens, illustrated in Figure 12, consists of a set of shaped coils which produce a toroidal field that deflects ions towards the axis. To first order, there is no magnetic field outside the coils, and the field inside is given

Figure 12. Geometry of an eight element toroidal field lens.

by $B(r,y) = B_{\theta 0} r_0 / r$, where r_0 is the inner radius of the lens and $B_{\theta 0}$ is the magnetic field at that position. The focusing properties are thus determined solely by the field boundary. Since the toroidal field lens has current elements within the beam volume, it is not subject to the off-axis aberrations of conventional magnetic or electric lenses. Another attractive feature is that space charge neutralization is easily accomplished by locating electron sources on the plates that hold the coils (Mohri, 1977), shown in Figure 12. Current cancellation of pulsed beams can also result from electron flow along the plates. The main disadvantage is that some of the beam is intercepted by the coils.

Coil boundaries can readily be determined for mono-energetic ion beams by numerical integration of orbits. In the approximation that only applied magnetic fields act on the ions, the normalized equations of motion are

$$\frac{d^2 R}{d\tau^2} = - \left(\frac{r_0}{r_{gi}} \right) \frac{V_z}{R} \quad (13)$$

$$V_z = \sqrt{1 - \left(\frac{dR}{d\tau} \right)^2}$$

where $\tau = \omega_{gi} t = (eB_{\theta 0} / m_i c) t$ and distances are scaled in terms of the ion gyroradius, $r_{gi} = m_i v_i c / eB_{\theta 0}$. A plot of typical ranges of r_{gi} for different ion species is given in Figure 13(a). If the incoming ion distribution and entrance field boundary are given, sample orbits are integrated until the condition $V_r / V_z = - R / (Z - Z_f)$ is satisfied, where Z_f is the desired focal point. This determines points on the exit boundary. The process can be inverted by integrating orbits back from a focus through a given exit boundary. Points on the

a)

Figure 13. Toroidal field lens calculations. (a) Scaled ion gyroradius versus beam energy and applied magnetic field. (b) Scale drawing of ion orbits through a toroidal field lens. Coil boundary calculated for 10 MeV He⁺ ions with a field at the inner radius of 25 kG.

entrance boundary are defined when sample ions are parallel to the axis. An example of such a calculation is shown in Figure 13(b). Parameters are appropriate for a 10 MeV He^+ beam. The 60 cm focal length requires a maximum field of 25 kG. For long focal length lenses, the exit boundary assumes a parabolic shape. The deflection must increase as r , while the field magnitude falls off as $1/r$, so the pathlength through the lens must increase as r^2 . We have developed a forward-reverse orbit integration scheme to calculate the exit and entrance boundaries of achromatic toroidal field lenses consisting of one or more elements. As expected, there must be a crossing of inner and outer particle orbits within the lenses. The boundaries must have a strong focusing effect on the ions. This means they are orientated at an oblique angle to the beam, and an achromatic focus can be achieved only over a narrow annulus. Although possible in principle, achromatic toroidal field lenses do not appear to be practical. Thus, energy spread will set one limit on the smallest focal spot attainable.

Limitations on ion beam focusing, both from processes within the lens and during transport, will now be considered. The effect of energy spread can be estimated by reference to Figure 14(a). The total deflection (α) is approximately proportional to $1/v_i$, so that $\Delta\alpha/\alpha \cong -\Delta E_i/2E_i$. This expression can be modified to give the aiming error

$$\Delta_f \cong \frac{1}{2} \left(\frac{\Delta E_i}{E_i} \right) R_\ell \quad (14)$$

b)

Figure 13(b).

Figure 14. Focal spot limitations. (a) Effect of a beam energy spread. (b) Effect of errors in locating field boundary. (c) Effect of initial beam divergence.

As an example, if a spot size of diameter 0.5 cm were required with $R_\ell = 15$ cm, then a total energy spread of 7 per cent could be tolerated. A second possible limiting factor is the accuracy of determining the field boundary, illustrated in Figure 14(b). If the beam has a total pathlength through the lens of s and the boundary is moved a distance Δs , then the aiming error is

$$\Delta f \cong \left(\frac{\Delta s}{s} \right) R_\ell \quad (15)$$

For a 0.5 cm diameter focal spot with $R_\ell = 15$ cm, the boundary need be accurate to only 1 cm if the total pathlength is 60 cm. Related errors can result from bulging of field lines between coils. Since, to first order, the ion deflection is proportional to the total flux lines traversed, the discrete coil errors can be made small. Another source of focal spot spreading is the inherent divergence (or lack of parallelism) of beams entering the lens. As shown in Figure 14(c), the divergence is usually categorized by a suitably averaged half-angle, $\Delta\theta_d$. In this case

$$\Delta f = \Delta\theta_d f \quad (16)$$

The distance to the target has been approximated as the focal length, f .

Since the electrons can not flow longitudinally in transverse magnetic field lenses, the self-magnetic field of high current beams can result in deflections of the ions. If the variation of beam current density is known, the self-field can

be calculated and the lens boundary adjusted to compensate for the additional deflection. This is not possible for rapidly pulsed beams with large variations in current. The effects of self-fields are reduced in toroidal field lenses by the presence of the coil plates (see Figure 12). Electrons accompanying the incoming ion beam are collected on the plates, flow through them, and then are emitted at the exit to neutralize the beam. Thus, the beam is divided into a number of smaller beams with isolated self-magnetic fields. This is the magnetic analogy of the multi-aperture injector for reducing electrostatic effects (Section II,A,1). If a beam with outer radius R_o and inner radius R_i is divided into N segments, then the ratio of the maximum field sampled by the ions to the maximum toroidal field of the full beam is

$$\frac{B_{\max}}{B_{\theta\max}} \sim \left(\frac{2\pi}{N}\right) \left(1 - \left[\frac{R_i}{R_o}\right]^2\right) \quad (17)$$

As an example, a reduction of self-field deflection by a factor of 2-3 can be accomplished in a 20 segment toroidal field lens.

To this point, we have considered focal spot limitations associated with the properties of the focusing element. Processes during subsequent propagation can also be important. Vacuum propagation will be discussed; propagation of high current ion beams in plasmas has been treated by a number of authors (see, for instance, Olson, 1980). Limits for completely unneutralized beams can be estimated using Equation (5). This equation holds for a cylindrical beam with uniform density and

no initial divergence. The first integral gives the following expression for the diameter of the minimum beam waist (the point where $dr/dz = 0$ for the envelope):

$$\frac{r_{\min}}{r_0} = \exp\left(\frac{-m_i v_i^3 \tan^2 \alpha}{4eI_B}\right) \quad (18)$$

In this expression, α is the maximum injection angle (the angle of the beam envelope as it enters the vacuum region), v_i is the ion velocity, I_B is the total beam current, and r_0 is the initial beam radius. Values of $r_{\min}/r_0$ are plotted in Figure 15 with $v_i = 2.5 \times 10^9$ cm/s (a typical value for inertial fusion applications) for H^+ , N^+ and Xe^+ . A radius decrease of a factor of 60 would be useful for inertial fusion. It is apparent that the focusing requirement severely limits beam current. High power can be achieved only for the heaviest ions at high energy. This is one of the prime motivations for Heavy Ion Fusion, in which it is proposed to use U^+ ions at an energy of 5-10 GeV at the 1 kA level (Martin, 1975; Maschke, 1975; Arnold, 1978).

Auto-neutralization allows much higher beam currents, but introduces new limitations on focusing caused by the presence of the electrons. If the electrons are injected with a transverse temperature, thermo-electric fields can be set up which defocus the ions (Humphries, 1978c). In other words, a term representing the electron pressure must be included in the envelope equation. This effect is small in a uniform radius beam, but can become significant in a focused beam where the electron distribution is compressed and heated. To estimate

Figure 15. Focal spot limitations for unneutralized and neutralized ion beams in terms of beam current in amperes. Solid lines: Unneutralized beams. Velocity equals 2.5×10^9 cm/s. Ratio of minimum beam waist to initial radius shown. Cylindrical beams with uniform density and zero emittance. Dashed lines: Focusing limits for neutralized beams set by electron pressure with initial electron temperature given by $kT_{e0} = 0.4 (m_e/m_i) m_i v_i^2$. Arrow indicates typical focusing requirement for inertial fusion applications.

this effect, it is necessary first to consider the distribution of transverse electric field in a neutralized beam with hot electrons. In a cylindrical beam with ion density specified by $n_i(r) = n_i(0)f(r)$ and uniform temperature electrons with density $n_e(r) \sim \exp(-e\phi(r)/KT_e)$, Poisson's equation can be written as

$$\frac{1}{R} \frac{\partial}{\partial R} R \frac{\partial \Phi}{\partial R} = e^{-\Phi} - \left(\frac{n_{i0}}{n_{e0}} \right) f(R) \quad (19)$$

where $\Phi = e\phi/kT_e$, $R = r/\lambda_d$, and $\lambda_d = (kT_e/4\pi e^2 n_{e0})^{1/2}$. Usually, the Debye length, λ_d , is small compared to the beam dimension; typical values are given in Part D of this section. Solutions to Equation (19) are plotted in Figure 16 for sharp boundary and parabolic ion density profiles. The electric fields do not rise linearly away from the axis; for the sharp boundary case, they are localized with a distance λ_d of the boundary. Thus, an exact envelope equation cannot be formulated. Detailed studies of radially dependent defocusing have been carried out by Lemons, Thode and Guiragossian (Guiragossian, et. al., 1980). A rough estimate can be derived if it is assumed that the ion density varies smoothly (as in Figure 16(b)) and the average electric field acting on ions is $E_r = \zeta kT_e/er$, where r is the envelope radius and ζ is a geometric correction factor. The electron temperature can be estimated in terms of the initial temperature by familiar reversible compression laws since focusing usually takes place over many electron bounce times. For a transverse

a)

Figure 16. Distributions of electron density and electric fields in cylindrical neutralized ion beams with uniform electron temperature. (a) Uniform ion density with sharp cut-off. (b) Parabolic ion density profile.

b)

Figure 16(b).

focus where there is enough longitudinal asymmetry to distribute electron energy in three dimensions, the following envelope equation holds (Humphries, et. al., 1981b)

$$\frac{d^2 r}{dz^2} \approx \frac{\zeta}{m v_i^2} \left(\frac{k T_{eo}}{r} \right) \left(\frac{r_o}{r} \right)^{4/3} \quad (20)$$

This equation determines the minimum beam radius as

$$\frac{r_{min}}{r_o} = \left[\frac{3\zeta}{4} \left(\frac{k T_{eo}}{1/2 m_i v_i^2} \right) \frac{1}{\tan^2 \alpha} \right]^{3/4} \quad (21)$$

which involves the injection angle, α , and initial electron temperature, $k T_{eo}$, but is independent of the ion current. To complete the argument, an expression for $k T_{eo}$ is needed. If the beam has been neutralized by longitudinal electron injection (Section III,B), then there must be a longitudinal electron temperature. If this is coupled into transverse electron energy, the computer simulations imply $k T_{eo} \sim .04 (m_e/m_i) (m_i v_i^2)$. Focal spot limitations for neutralized beams have been plotted in Figure 15 using this expression. At high current levels, a great advantage can be gained by neutralization. Even light ions (H^+) are close to meeting the fusion focusing requirement. For heavier ions, electron thermal effects do not constitute a significant problem.

D. Ion Beam Transport by Electron Control

Intense ion beams can generate such strong potentials that they cannot propagate without the addition of electrons. This

fact can be used to advantage. By controlling the location of electrons, ions can be directed (Humphries and Poukey, 1979). In this bootstrapping process, electrons can be influenced by magnetic fields much too weak to confine low velocity ions; electric fields resulting from charge separation can then act effectively on the ions. Experiments on this effect have been reported by Robertson (Robertson, 1981).

An idealized electron control transport system is shown in Figure 17(a). A high current ion beam enters a region with an applied solenoidal magnetic field through a localized electron source. The boundary condition at this source ($E_z = 0$) implies that equal numbers of ions and electrons enter the volume. The transport system must counteract the effect of the ion beam emittance, which can be characterized by a transverse beam temperature, kT_i . Downstream, the ion distribution tries to expand while the electrons are confined by the magnetic field. Ion expansion leaves an excess of electrons in the original column, so a charge separation sheath is set up which reflects the ions. Exact sheath solutions (which depend on the ion transverse velocity distribution) have not been calculated, but sheath dimensions can be estimated from some simple scaling laws. If electrons are strongly confined by the magnetic field, the sheath will be similar to Figure 16(a) with the roles of the ions and electrons reversed. The sheath width is on the order of the Debye length calculated in terms of the ion temperature,

a)

Figure 17. Ion transport by electron confinement. (a) Idealized geometry with solenoidal field. (b) Scaling of electrostatic and magnetic sheath widths as a function of electron and/or ion temperature. For electrostatic sheath, use density scale; for magnetic sheath, use magnetic field scale.

$\lambda_E = \sqrt{kT_i/4\pi e^2 n_i}$. For weak magnetic fields, the sheath width is determined by the magnetic confinement of the electrons.

As discussed in Section IV, the sheath width for magnetic confinement is approximately $\lambda_M = \sqrt{2ekT_i m_e c^2 / e^2 B_0^2}$. A combined plot of electrostatic and magnetic sheath widths is given in Figure 17(b). The combined sheath width is on the order of the maximum allowed value. The plot indicates that even with modest magnetic fields the sheath width is about λ_E . For high current beams, λ_E is small; at low density the sheath approaches centimeter size, so the method is not suitable for low current beams.

Two examples of practical electron control transport systems are shown in Figure 18. The first utilizes solenoidal fields with electron sources connected to the beam through fringing field lines. For pulsed beams with length shorter than the transport system, electrons join the beam at the entrance by a combination of transverse and longitudinal neutralization processes and travel with the beam to provide current cancellation. The ion beam must give up a fraction (m_e/m_i) of its energy to accelerate the electrons. For a long pulse beam, the beam potential adjusts so that there is a thermal distribution of electrons with zero net longitudinal velocity. The time over which transport will be effective depends on the stability of the electron distribution against cross field motion. Experiments have indicated that electrons are generally unstable in this situation (Maron, 1981; Humphries and Kuswa, 1979; Humphries, et. al., 1980a), but transverse diffusion can

b)

Figure 17(b).

a)

Figure 18. Practical geometries for ion transport by neutralizing electron confinement. (a) Modified solenoidal field. (b) Multiple cusp fields.

b)

Figure 18(b).

be slow. In this case, it may be possible to transport steady-state beams by incorporating sweeping electric fields to remove diffusing electrons. Plates at a positive voltage with respect to the source could intercept magnetic field lines external to the beam volume. The second geometry, which uses multipole magnetic fields to exclude electrons, leaves the beam volume almost field free. Again, sweeping fields could be used to prevent the growth of an electron density in the magnetic field regions.

IV. HIGH CURRENT ION BEAM ACCELERATOR THEORY

Vacuum auto-neutralization allows the transport of high current ion beams. We now turn attention to the production of such beams. The major problem in neutralized beam accelerators is controlling the electrons in the presence of applied electric fields. If conventional gaps were used, there would be a reverse electron flow which would waste power and prevent effective neutralization in the intervening drift regions. Drift regions must be isolated from acceleration gaps. This can be accomplished by grids or foils, but the presence of conductors in the beam path is impractical for high power accelerators. Studies in the Pulselac Program have concentrated on applied magnetic fields to control electrons. The first part of this section concerns the physics of magnetically insulated acceleration gaps. The second part deals with system problems that can arise when such gaps are operated in series.

A. Magnetically Insulated Acceleration Gaps

A schematic diagram of a magnetically insulated acceleration gap (Sudan and Lovelace, 1973) is shown in Figure 19. The transverse magnetic field applied is strong enough to prevent electrons from crossing the gap. Ions, being more massive, pass through with small deflection. The critical magnetic field is the minimum field necessary for electron insulation and is given by

$$B_{\text{crit}} = \sqrt{\left(\frac{2eV_0 m_e c^2}{e^2 d^2}\right) \left(1 + \frac{eV_0}{2m_e c^2}\right)} \quad (22)$$

Figure 19. Idealized one-dimensional, magnetically insulated, ion injector.

For higher fields, the penetration depth of electrons into the gap, Δ , decreases approximately as $(\Delta/d) \sim (B_{\text{crit}}/B)^2$. Insulating field requirements are modest. For a 1 MV applied voltage on a 1 cm gap, $B_{\text{crit}} = 4.8$ kG. The ability to apply high electric field stress to gaps is essential for high current beams. Since the beams are almost unneutralized in the gaps, the fraction of the accelerator length occupied by gaps must be minimized.

The geometry of a practical magnetically insulated gap is more complicated than that of Figure 19. Electrons have a drift motion perpendicular to $\vec{E}$ and $\vec{B}$. The gap geometry must be symmetric along this direction to prevent bunching of electrons (Humphries, et. al., 1977). Electron bunching results in electric field components along the drift direction which can induce electron loss across the gap. If the accelerating field is longitudinal (E_z), the two choices for applied magnetic field are B_θ and B_r . Besides presenting practical difficulties, a toroidal field causes radial electron drift with a clear violation of symmetry. Radial fields, on the other hand, result in an azimuthal drift of electrons which can be self-connected. It is thus possible to have an equilibrium of confined electrons. The radial field gap (Humphries, 1978b) which has precedents in the field of plasma accelerators (Morozov, et. al., 1972; Cole, et. al., 1974) is illustrated in Figure 20(a). It has additional attractive features for neutralized ion beam acceleration. Magnetic deflections of ions traversing a gap are predominantly azimuthal; since the number of inwardly and outwardly directed

Figure 20. Radial magnetic field ion acceleration gap. (a) Three-dimensional view. (b) Cross-section showing magnetic field lines.

flux lines are equal, there is no net change in ion direction. The radial field gap is one example of a zero canonical angular momentum geometry. Flux lines near the gap connect directly to conducting surfaces which can act as electron sources, facilitating transverse neutralization. The main complicating factor associated with radial magnetic fields is that the origin (where $B_r = 0$) cannot be used for beam transport. Thus, the beam must have an annular cross-section, as shown.

The radial field gap is well suited to the formation and use of virtual gap electrodes. The existence of virtual cathodes has been well documented in intense ion diode research (see, for instance, Humphries, 1980a). In the presence of applied electric fields, electrons emitted at the cathode flow along transverse field lines to provide neutralization. The magnetic fields thus determine equipotential surfaces. The electron cloud acts like a conducting foil of high transparency; the part of the electron cloud adjacent to the accelerating gap is usually called the virtual cathode. Image current neutralization (see Section II,A,1) can occur in a virtual cathode, so the cathode extraction aperture can be much larger than the gap spacing. Similar considerations apply to the anode side of the gap, although the boundary of the electron cloud need not be coincident with the magnetic field lines. The anode equipotential occurs in a plane determined by a balance between the attraction of the bare beam space charge in the gap and the repulsion of the applied field. The location of virtual electrodes is indicated in Figure 20(b). With electrons present, the

ρ_1

shape of the accelerating fields are determined by the shape of the magnetic fields. The transverse components of the electric fields can then be used to provide effective electrostatic focusing of the ions. (This does not preclude longitudinal focusing, since the well-known arguments about the interdependence of longitudinal and transverse stability (see, for instance, Vlasov, 1968) are not valid in the presence of electron space charge.) If the virtual electrodes have an average radius of curvature ρ_l (as shown in Figure 20(b)), then in the continuous approximation, the effective focusing potential of a series of gaps separated by a distance l is

$$U_l \approx (eV_0) \left(\frac{X^2}{2\rho_l l} \right) \quad (23)$$

where V_0 is the applied potential and X is the distance from the average beam radius. The depth of the confining potential is independent of the ion mass and energy. Thus, if the beam emittance is preserved, electrostatic focusing can be effective throughout the acceleration process.

The shape of the magnetic field in the gap is important for transverse focusing. The field patterns shown in Figure 20(b) were calculated using a computer code with the capability of handling pulsed fields in the presence of imperfect conductors (Freeman, 1981). There are a variety of methods to change the field curvature in actual gaps. If the field has a variation of curvature with position, the location of virtual electrodes can be moved by

the use of adjustable extensions on the cathode and anode structures. At a fixed location, the curvature can be varied by changes in the relative currents in the upstream and downstream coil sets. For rapidly pulsed fields, a conducting plate near the gap centerline can be moved longitudinally to change the average curvature.

Electron and ion behavior in a radial field gap have been studied with two-dimensional particle simulation models (Poukey, 1979). The results shown in Figure 21 were generated to investigate downstream electron behavior. A fixed proton beam of 10 kA with an initial energy of 1 MeV traversed the 0.2 MV gap. Electrons were emitted from the boundaries shown with a flux determined by space charge limitations. The calculation was continued until a steady-state was reached. Plots of the applied equipotentials and magnetic field lines are given in Figure 21(a) and (b). Equipotentials with the inclusion of electron and ion space charge at 27 ns are shown in Figure 21(c). There is almost complete neutralization in regions connected to the cathode structure and a well-defined virtual cathode oriented along a magnetic field line. The only factor acting to randomize the electron distribution was the two-dimensional geometry of the problem. A configuration space map of the electron distribution is given in Figure 21(d). A high magnetic field ($\sim 5B_{\text{crit}}$) was used in this run to demonstrate clearly the virtual cathode. At lower field, electrons penetrated further into the gap so that the virtual cathode was not a sharply defined surface, but there was still good neutralization in downstream regions.

Figure 21. Two-dimensional computer simulations of radial magnetic field gap for ion beam post-acceleration. Gap width: 0.7 cm. Gap voltage: 0.2 MV. Average applied magnetic field: $5 B_{crit}$. Proton beam (10 kA, 1 MeV) introduced from left-hand border. Space charge limited flux of electrons introduced on cathode surfaces. (a) Equipotential lines for applied electric field only. (b) Applied radial magnetic field lines. (c) Equipotential lines with the introduction of both electrons and ions. Note virtual cathode and downstream neutralization. (d) Electron distribution. Electron emission surfaces are indicated by dashed lines.

The injector gap is a special case of the magnetically insulated gap in which low velocity ions are produced on a real anode. The study of injector gaps, with their attendant longitudinal space charge limits, underlies the field of high current ion diodes. The physics of injectors is one of the most difficult problems for neutralized ion accelerators. Self-consistent theories have been developed for ion flow in magnetic gaps with cathode emitted electrons (Sudan and Lovelace, 1973; Bergeron, 1976; Antonsen and Ott, 1976). These equilibrium models show that at practical magnetic field strength ($B > 1.5 B_{crit}$), electrons have little effect on ion flow, which is space charge limited to about j_{iC} . As an example, j_{iC} for N^+ ions in a 0.5 cm gap with 100 kV applied voltage is less than 2 A/cm^2 . Recent theoretical studies (Poukey, et al., 1981a), motivated by observations of ion flow from injectors more than an order of magnitude above the Child limit (Humphries and Kuswa, 1979; Humphries, et al., 1980a; Greenly, 1980), have indicated how enhanced ion flow can occur in magnetically insulated gaps. In the radial field gap, there are a variety of possible trapped electron orbits. Enhanced ion flow can occur if the trapped electrons are not constricted near the cathode region, but are allowed to fill the gap. In a sense, the ion beam is partially neutralized within the gap. Modified electron distributions can result if the gap is prefilled with electrons, or if there is a mild diffusion of

electrons across field lines. Electron prefill can be accomplished in practice by filling the gap with a plasma before application of the injector voltage. Following the rise in voltage, there is a phase during which plasma ions are cleared from the gap but plasma electrons are trapped. These electrons can neutralize ions emitted from the anode, which flow across the gap with a density about equal to the initial plasma density. The results of one-dimensional computer simulations of this process (Poukey, et al., 1981a) (including plasma particles and space charge emitted ions and electrons) are shown in Figure 22. The following parameters were used: 1 cm gap spacing, 1 MV applied voltage, 9.6 kG applied field ($2B_{crit}$), $1.2 \times 10^{13} \text{ cm}^{-3}$ plasma prefill density, and H^+ ions. The voltage rose to its steady state value in 5 ns, comparable to the ion clearing time, and long compared to $1/\omega_{pe}$. In agreement with experiments, the ion flow was significantly enhanced immediately following the clearing phase (Figure 22(a)). During the voltage rise, all the prefill electrons remained trapped in the gap. The spatial variation of potential in the final state (10 ns) is shown in Figure 22(b). Most of the potential drop occurred near the anode, accounting for the enhanced flow. Electron and ion densities are plotted in Figure 22(c); the electron density actually increased toward the anode. A plot of magnetic flux density versus position is shown in Figure 22(d). For the relativistic gap treated, diamagnetic effects of the electron drift motion can be important, particularly since the total number of trapped electrons is about a factor of 40 higher than

Figure 22. Computer simulations of ion flow enhancement in a one-dimensional, magnetically insulated, ion injector gap with plasma prefill. Electrons and protons emitted from cathode and anode, respectively, with $E_x = 0$ boundary condition. Initial plasma electron and ion densities: $1.2 \times 10^{13} \text{ cm}^{-3}$. Gap spacing: 1 cm. Final applied voltage: 1 MV. Voltage risetime: 4 ns. Applied magnetic field: 9.6 KG ($2 B_{\text{crit}}$). (a) Ion flow enhancement (solid line) and applied voltage (dashed line) as a function of time. (b) Spatial variation of potential at 8 ns. (c) Electron and ion densities at 8 ns. Thin dashed line represents initial plasma density. (d) Spatial variation of magnetic field at 8 ns. Thin dashed line indicates vacuum field.

Figure 22 (b).

Figure 22(c).

Figure 22 (d).

a space charge limited gap. Flux lines were compressed toward the anode, contributing to the flow enhancement. For the fast pulse considered, total flux was conserved between the two conducting electrodes, so that the gap remained insulated despite the diamagnetic electron currents.

B. Multi-Gap Effects in High Current Ion Accelerators

In addition to raising particles to high energy, an accelerator must also maintain qualities of the particle distribution so that the beam can be transported or optimally focused after exiting. If there are no collisions, Liouville's Theorem implies that the total phase volume of the beam distribution $f(x, v_x, y, v_y, z, v_z)$ is a constant during acceleration (see, for instance, Lawson, 1977). Even though the volume remains constant, in some circumstances the phase distribution can become filamented. The effective volume occupied by the diffuse distribution determines transport properties. In this section, we will consider processes occurring in high current, multi-gap systems that can influence the effective volume of the beam distribution. To first order, transverse and longitudinal beam motions are decoupled, so it is possible to consider these problems separately. In this paper, we will refer directly to phase space volume. Most papers on accelerators use the empirically motivated quantity of emittance, where particle orbits are characterized by position and angle $(x-dx/dz)$ rather than velocity $(x-v_x)$. In terms of these variables, the area inside a suitable defined boundary is

called the emittance, ϵ_x . This quantity is not conserved during acceleration although another quantity, the normalized emittance, is. The normalized emittance, given by $\epsilon_{xn} = (v_i/c)\epsilon_x$, is equivalent to the area of the phase space distribution.

1. Transverse Beam Control

In a multi-lens transport system, there is a maximum transverse beam dimension. This, combined with the strength of the lenses, defines an upper limit on transverse velocity that can be contained (see Section II,2). Thus, for any transport channel, there is a maximum transverse phase volume of particles that can be accepted. If a beam has a sufficiently small phase volume, its shape can be adjusted by a combination of focusing elements to fit into the transport system. The benefits of minimizing transverse phase volume for beam focusing can be seen in Figure 23. The beam distributions entering the lens and at the focal point are indicated. In the limit that the maximum focusing angle (α) is large compared to the beam divergence ($\Delta\theta_d$), conservation of phase volume yields an expression for the minimum focal spot size, $\Delta_f \geq \Delta\theta_d R_l/\alpha \simeq \Delta\theta_d f$. Extension to two transverse dimensions is straightforward. The power density that can be delivered at a focus is proportional to $B = I_b/\epsilon_x\epsilon_y \simeq j_i/\Delta\theta_d^2$, where I_b is the beam current and j_i is the current density at the lens. The quantity B is called the brightness.

In low current accelerators (which form the basis for most experience to date), space charge forces are small and applied

Figure 23. Calculation of focusing limits set by emittance of beam entering lens.

focusing fields can be set up which vary linearly away from the beam symmetry axis. In this case, it is well known that beams can be propagated without transverse phase space filamentation. An example for a slit beam in a cylindrical lens array is shown in Figure 24(a). In this case, there are few lenses per phase oscillation (large phase advance per cell); the elliptical phase distribution is shown at the entrance to five sequential lenses. Even though the orientation of the distribution is changed, it maintains its continuity. Recent work on high current, space charge dominated ion beams has shown that significant growth in the effective phase volume can occur in a lens array (Hofmann, 1981; Haber, 1979; Laslett and Smith, 1979; Penner and Galejs, 1979). This comes about because of non-linear forces contributed by the beam self-fields. Net forces can be highly non-linear in neutralized ion beams. Two examples have been shown: the image charge neutralized acceleration gap (Section III) and the thermo-electric fields of a neutralized beam (Figure 16). The effect of introducing non-linear components into the lens deflections can be seen in Figure 24(b). After 50 lenses, the distribution is filamented and occupies a larger effective phase volume.

The emittance growth in Figure 24(b) is not a result of an instability, since the lens parameters have been chosen so that all particle orbits are stable and there are no self-field interactions. The cause can be seen by consideration of Figure 24(a). In this case (with few lenses per phase oscillation), the distribution assumes skewed orientations with respect to

Figure 24. Phase space plots for slab beams traversing an array of lenses with uniform spacing, ℓ . Initial elliptical distributions appear as circles because of plotting normalization. (a) Linear lenses with focal length, f_0 , equal to 1.1ℓ (large phase advance). Initial distribution and distributions following the first 5 lenses are plotted. (b) Lenses with large phase advance ($f_0 = 1.1 \ell$) and nonlinearities of the form $f(x) = f_0(1 - 0.1(x/x_0)^2)$, where x_0 is the initial width of the distribution. Initial distribution and distributions following transport through 50 lenses is shown. Note filamentation of distribution. (c) Linear lenses with small phase advance ($f_0 = 20 \ell$). Initial distribution and distributions following first 5 lenses. Note that distribution is approximately stationary. (d) Lenses with small phase advance ($f_0 = 20 \ell$) and same non-linearity as in (b). Initial distribution and distributions following 225 lenses (approximately same number of phase oscillations as in (b)) are plotted. Note small emittance growth.

the axes so that it sweeps across a larger phase area than it occupies. The distribution does not grow in the linear case because all particles rotate in phase space at the same oscillation frequency. In the non-linear case, there is a variation of oscillation frequency, so that eventually phase mixing will spread the distribution throughout the region swept out. Once this occurs, emittance growth stops. A means of solving this problem is indicated in Figure 24(c). Here, the focusing strength of the linear lens array ($1/f\ell$ in the nomenclature of Section II,2) has been maintained constant, but the number of lenses per phase oscillation has been increased. Orbits in the lens array approach the continuous approximation. We have also chosen a beam distribution with a boundary corresponding to the phase space orbit of a peripheral particle. In this case, the phase distribution maintains an almost constant orientation. The phase space area swept out by the beam is only slightly larger than the area of the distribution. Thus, there is no emittance growth induced by variation of phase space rotation frequencies. This conclusion is confirmed in Figure 24(d), where the same relative non-linearity has been added to the lens force as in Figure 24(b). The distributions entering the first lens and after 250 lenses are shown. We state a straightforward but important conclusion: approximately stationary phase distributions can be propagated in non-linear focusing systems in the continuous limit (small phase advance per cell). Typical parameters for the electrostatic focusing method described in Section IV,A put it well

into the continuous limit. For the strong focusing quadrupole systems typical of most conventional accelerators, it is not practical to use small phase advance. This is because the net focusing effect of a quadrupole set scales as the cube of the length of the lens (Lawson, 1977) rather than linearly. Thus, if the number of lenses per length were increased to approach the continuous limit, the applied fields would have to be increased to achieve the same focusing strength. It is apparent that phase mixing and emittance growth are an inevitable problem in practical alternating gradient focusing systems if space charge becomes dominant. For this reason, we will concentrate attention on first order, electrostatic focusing of high current beams in the continuous lens limit.

The transverse dynamics of high current beams can be treated easily using the average potential energy model (Humphries and Poukey, 1979). Three major effects will be considered: (a) electrostatic focusing in acceleration gaps, (b) space charge deflections from bare ion space charge in gaps, and (c) toroidal self-magnetic field deflections from lack of current neutralization in the transverse magnetic fields near the gaps. It is assumed that in regions far from the gaps, the beams are neutralized and travel ballistically. The gap focusing has already been treated (Equation 23). Space charge defocusing can be estimated from the calculations of Section II,1 for uniform beam density. Independent of gap width, the maximum electrostatic potential at the edge of the beam is $0.5 \phi_0$. Averaging the edge potential across the gap

reduces this value by a factor 0.7; an additional factor (d/ℓ) accounts for the fact that defocusing occurs only in the gap. The maximum effective potential energy for self-magnetic fields acting on a thin annular beam is

$$U_t(\text{max}) = 2\pi e^2 n_{i0} \left(\frac{v_i}{c} \right) \left(\frac{d^*}{\ell} \right) (\chi + \Delta\chi)^2$$

where it is assumed that beam current is uncanceled in a region d^* near the gap. Typical variations of the effective potential energy and the addition of all effects is indicated in Figure 25. The parameters are appropriate to a 15 kA N^+ beam with $j_i = 50 \text{ A/cm}^2$, average radius 10 cm, gap voltage 1 MV, and $\rho_\ell = 20 \text{ cm}$. A plot is shown for beam energy of 1 MeV. The self-magnetic field potential energy is not symmetric about the average beam radius. Inspection of Figure 25 implies the following conclusions. (1) the net potential energy well has a skewed shape that is not consistent with the assumption of uniform beam density. Calculations of self-consistent beam distributions have been described (Humphries and Poukey, 1979). (2) In principle, an ideally matched cold beam could be injected into the transport system shown and propagated without emittance growth. In practice, the injected beam will be non-ideal and must adjust itself to the transport system, developing a transverse temperature. The resulting beam temperature will probably be a significant fraction of the defocusing potential energy. (3) Space charge defocusing is most severe at the accelerator entrance, and decreases downstream with decreasing beam density. The self-magnetic field effects increase with the beam velocity. The

Figure 25. Examples of average potential calculations for transport of ion beams with strong self-field effects. U_f : Electrostatic acceleration gap focusing potential. U_e : Space charge effect of bare beam in acceleration gaps. U_m : Self-magnetic field effects. U_T : Total confinement potential. 15 kA N⁺ beam, 1 MeV, 1 cm gap width, 1 m between gaps, self-magnetic fields act over 10 percent of the intergap distance, $\rho_l = 20$ cm, average beam radius 10 cm, $j_i = 50$ A/cm².

axi-centered self-magnetic field focusing could be counteracted by offsetting the gap coil sets and shaping the virtual electrodes (Humphries, 1977). There is a limit to how accurately this can be done, so that transverse beam temperature of 100 eV or higher is likely. This does not preclude applications to inertial confinement fusion; a divergence limited focus of a 100 MeV N^+ beam to a spot 0.5 cm in diameter 2 m from the lens requires a beam divergence of 1.3 mrad, equivalent to a 200 eV transverse temperature.

2. Longitudinal Beam Control

Transport through and beyond accelerators also places constraints on the longitudinal phase volume of beams. In radio-frequency accelerators, particles must remain bunched on the stable portion of the accelerating waveform. In pulsed accelerators (such as linear induction accelerators (Christofilos and Hester, 1964)) the beam passage time must be limited to the pulselength of the gap fields and the beam current must provide proper gap loading. Longitudinal phase dynamics have not been of great concern in existing linear induction accelerators (Leiss, 1979) which produce relativistic electron beams. Applications to inertial fusion have prompted interest in induction accelerators for non-relativistic ion beams (Faltens, et al., 1979; Humphries, et al., 1981a). For ions, longitudinal dynamics may be the most severe problem area. A variety of approaches has been taken towards the investigation of beam confinement and stability. We have

developed a velocity lens model (Humphries, 1981), in which the motion of particles about the center of mass of the beam pulse can be shown to approximate the transverse motions of particles through a series of lenses. The voltage levels in the gaps act as velocity focusing forces. A voltage waveform which rises during the passage of the beam acts as a positive focal length lens. Slower particles which lag behind are accelerated more so that they catch up with the rest of the beam, and the inverse holds true for faster particles. The linear voltage ramp shown in Figure 26(a) is the equivalent of a linear lens. Familiar results from lens array theory (see Section II,1) can be used to find the system acceptance in longitudinal phase space, to calculate beam bunching, and to investigate microinstabilities. The latter can occur when the beam passes through an active focusing medium with the index of refraction a function of the beam fluence (the effect of current on the voltage waveform). A waveform suggested by Faltens (Faltens, 1980) for the confinement of a beam with a flat-top in current is shown in Figure 26(b). This acts as a non-linear lens (with an approximately square effective potential well). If there are many acceleration gaps per longitudinal phase oscillation (the continuous limit), it has been shown in Section IV,B,1 that a beam can be propagated with little increase in emittance.

The cavity of an inductive linear accelerator (see Section V,D) has a low Q , so there must be significant loading of the gap voltages by the beam current for efficiency. In this case,

Figure 26. Examples of voltage waveforms in acceleration gaps to provide longitudinal beam confinement. (a) Linear voltage ramp. (b) Synthesis of voltage waveform to provide non-linear longitudinal confinement of a beam with a flat-top in current.

velocity bunching instabilities of non-relativistic beams can occur. A current perturbation in a downstream gap results in a distortion of the longitudinal phase distribution through the gap voltage response. Velocity bunching can then occur, which may result in a larger current perturbation in a downstream gap. This process has been investigated for neutralized beams of uniform velocity passing through a periodic series of gaps (Humphries, 1980b). The equivalent circuit for the beam-gap system is shown in Figure 27(a). Only the AC component of voltage was preserved; reactive components of the gap impedance were neglected. A phase space plot of an unstable beam is shown in Figure 27(b). Note the phase space distortion that results in a downstream current spike. The condition for stability is

$$\Delta v / v_i \geq \mu / (\mu + 2) \quad (24)$$

where $\Delta v / v_i$ is the longitudinal velocity spread and μ is defined in Figure 27(a). Increasing circuit damping (lower μ) lowers the dependence of gap voltage on beam current and is stabilizing. Increasing the velocity spread decreases the relative bunching that occurs with a given voltage oscillation and also acts to stabilize the beam. Since the instability results in a larger longitudinal velocity spread, it is self-quenching. The maximum velocity spread is limited by the acceptance of the system. Large velocity spread limits longitudinal bunching and can also affect transverse focusing if the lens deflections depend on energy. The velocity lens

a)

Figure 27. Computer simulations of longitudinal instability of a non-relativistic ion beam passing through resistively loaded cavities. (a) Equivalent circuit for ideal linear inductive accelerator cavity driven by a pulse-forming-network. Steady state beam provides a load R_0 ; damping resistors with resistance μR_0 can be included to decouple cavity voltage from beam current. Simulation treats periodic beam with AC components of voltage only. (b) Typical simulation results for cavities with insufficient damping. Beam phase space plots in rest frame with z on abscissa and v_z on ordinate. Numbers indicate gap number. Plots are scaled; solid black line indicates initial velocity spread.

b)

Figure 27(b).

model has been used to study effects of reactive circuit elements in cavities. Longitudinal stability is presently an active area of research. Promising directions for reducing velocity bunching are the use of adequate cavity circuit damping, capacitive shunting of high frequency gap voltage oscillations, inclusion of microwave absorbers in the cavity to damp out resonances, and the use of active clamping elements to maintain a constant voltage under a changing beam load (Humphries, 1977).

While velocity bunching presents problems for non-relativistic beams, it can also be used to advantage to achieve maximum beam power for inertial fusion applications. Even a small longitudinal compression after leaving the accelerator would be useful. For instance, a bunching factor of 5 could reduce the energy delivery time for a beam from 50 ns (a lower value for practical inductive linear accelerator technology) to the 10 ns pulselength required for fusion drivers (see, for instance, Clauser, 1975). The longitudinal phase volume of a beam limits bunching just as the transverse phase volume limits focusing. If a beam with initial pulselength Δt_p and velocity spread $\Delta v/v_i$ passes through an ideal velocity lens with a voltage waveform which induces a linear variation in beam velocity, $\pm v_b$, then the minimum beam pulselength is

$$\Delta t_p(\text{min}) \simeq \left(\frac{\Delta v_z}{v_b} \right) \Delta t_p \quad (25)$$

High current beams with small longitudinal velocity spread passing through a series of inductively loaded cavities (as shown in Figure 28) can be unstable. This instability can induce self-bunching of a beam without the need for active bunching gap switching. The beam current across the cavity inductance generates the negative-positive bunching waveform. The results of computer simulations of this process are shown in Figure 28. Ideally, the beam density should vary longitudinally as $n_i(z)/n_{i0} = 1 - [(z-\bar{z})/\Delta z]^2$ where $\bar{z}$ is the average position and Δz is the beam length, to generate a linear voltage ramp. An elliptical phase distribution with $n_i(z)/n_{i0} = [1 - [(z-\bar{z})/\Delta z]^2]^{1/2}$ was chosen to demonstrate the effects of a non-ideal current waveform. A 10 MeV He^+ beam with velocity 2.2×10^9 cm/sec and a 1 percent velocity spread was used. It passed through four bunching gaps with voltage swings of 0.6 MeV. The beam current profile leaving the gaps, 2.8 meters downstream, and 6.6 meters downstream are shown. Bunching was clearly limited by velocity lens imperfections rather than the longitudinal velocity spread. A compression factor of 3.7 was achieved, an interesting figure for inertial fusion applications.

Figure 28. Method for longitudinally bunching non-relativistic beams using inductively loaded cavities. Current profiles generated by a computer simulation are shown at various positions from bunching gaps.

V. EXPERIMENTS ON INTENSE ION BEAM ACCELERATION

Experimental work related to intense ion beam generation using electron neutralization has been carried on for a number of years with the goal of developing high current ion diodes. A review of this field has recently been prepared (Humphries, 1980a). This section will be limited to a summary of experimental work directed toward high current accelerators in the Pulselac Program. (Humphries, et al., 1979a; Humphries, et al., 1979b; Humphries, et al., 1980b; Humphries, et al., 1981a.)

A. Ion Sources

Ion sources provide a high flux of low energy ions which can be formed into a beam in the injector gap. The flux available from a stationary plasma is determined by the Bohm limit (Bohm 1949).

$$j_i \approx 0.4 n_p \left(\frac{2kT_e}{m_i} \right)^{1/2} \quad (26)$$

where n_p is the plasma ion density. The transport of ions to the plasma surface is by diffusion. Thermal velocities are low for ions, so that a high flux source must have a high density plasma. Most work to date has been performed with capacitively ballasted flashboards, which utilize a large number of surface sparks to produce a dense plasma on the anode (Humphries, et al.; 1976, Dreike, et al., 1976). Flashboards are good sources of protons and other light ions for short pulse work. They have the disadvantage that they

permit no control of plasma flux and produce considerable background contamination.

In order to produce controlled fluxes of ions, we have concentrated on active plasma sources external to the injector gap (Humphries, et al., 1981d). High flux was achieved by imparting a directed velocity to the plasma by magnetic acceleration and directing it into the injector gap. The available flux in this case is given by

$$j_i \approx 0.5 n_p V_A \approx 0.5 n_p \sqrt{\frac{B_a^2}{4\pi m_i n_p}} \quad (27)$$

where B_a is the average pulsed magnetic field acting on the plasma and V_A is the Alfvén velocity. Directed ion energy in the range of 1 KeV can be obtained with pulsed fields of a few kG, increasing the flux more than an order of magnitude above the thermal limit.

Two plasma guns suitable for producing an annular plasma flow for high current accelerator applications are shown in Figure 29(a) and (b). The first used surface sparks on solid insulators to produce a hydro-carbon plasma prefill. The plasma was then accelerated longitudinally by the pulsed radial magnetic field produced by two concentric coils. Proton fluxes exceeding 100 A/cm^2 (equivalent) were produced for pulse lengths of about $0.5 \mu\text{s}$, with a useful extraction area of about 200 cm^2 . A second gun (Figure 29(b)) used a fast pulsed pancake coil with photoetched, spiral windings to produce a radial field in a localized gas layer. The gas was introduced through a fast

Figure 29. Plasma guns to provide high flux rings of longitudinally directed plasma for short pulse ($1 \mu s$) ion extraction. (a) Magnetic acceleration source using surfaces sparks to provide an initial low energy prefill. A: Metal plasma expansion chamber. B: Power inputs for pulsed magnetic field. C and D: Coaxial plotted magnet coils to provide a pulsed radial field. E: Spark source with sample holder (12 used). F: Power input for spark source. G: Extraction slots. (b) Gas injection source. A: Main support flange (lucite). B: Lucite plasma expansion chamber. C: Extraction slots in field excluder. D: Pulsed gas valve. E: Ceramic gas distribution nozzle. F: Pre-ionization electrode. G: Spiral coil, photoetched and plotted. H: Pyrex plate.

Figure 29(b).

valve. The pulsed field ionized the gas layer and accelerated it longitudinally with little included flux. This gun allowed the generation of a wide variety of heavier ion species. Purity levels of about 80 percent N^+ and Ne^+ have been obtained in μs pulses with flux exceeding $50 A/cm^2$ (equivalent).

B. Electron Sources

Neutralization depends on the availability of practical electron sources. In most experiments to date, sufficient electrons were produced by secondary emission from surrounding walls bombarded by part of the primary ion beam. The next generation of ion accelerators will need active electron sources to prevent excessive beam loss. Thermionic sources are impractical for pulsed beam applications. High electron flux is necessary for a small fraction of the time ($\sim 10^{-6}$). A new type of active cathode utilizing electron extraction from surface spark plasmas has been investigated at a number of laboratories (Ramirez and Cook, 1978; Humphries, 1978d; Loda, 1979; Fessenden, et al., 1981). Many sparks are necessary to produce large area plasmas, so ballasting is necessary. A capacitively ballasted plasma cathode (Humphries, et al., 1981e) is shown in Figure 30. It consists of a number of sealed layers. From the bottom upwards, these are an insulator, a high voltage electrode, an insulating dielectric layer, and a ground plane. The grounded electrode is photoetched with a series of holes. When the high voltage electrode is pulsed with an oscillating waveform at about

Figure 30. Capacitively-ballasted, surface spark electron source.

20 kV, there is high field stress (>100 kV/cm) where the ground electrode contacts the dielectric. Sparks are initiated, and plasma fills each of the holes. Energy is coupled to the plasma through the capacitance between the two electrodes. The source has fairly low available energy, and can supply about 10 A/cm^2 of electrons in a $1 \mu\text{s}$ pulse. Being thin and flexible, it can be mounted easily inside drift tube structures. Higher current density can be achieved with resistively ballasted flashboards.

C. Ion Injectors

The injector gap, which accelerates ions from low energy to 0.1 MeV or higher, presents the greatest difficulties in the system. It must operate in conjunction with the plasma source and must overcome severe space charge limitations. Considerable experimental work has been performed to investigate the methods of ion current density enhancement discussed in Section IV,A. Initial evidence for trapped electron enhancement of ion flow came from experiments on long pulse ($1 \mu\text{s}$) gaps with proton flashboards (Humphries and Kuswa, 1979). The ion flow was found to rise continuously over the pulse, ultimately reaching levels a factor of 30 above the Child limit. Net ion flux was azimuthally symmetric and the growth rate for enhancement was sensitive to the strength of the magnetic field, indicating electron diffusion as the controlling mechanism. The first indication of ion flow enhancement by electron prefill was found in experiments in which a 300 eV carbon plasma was injected into the gap (Humphries, et al., 1979c). There was

a clearing phase and subsequent enhanced flow at a level influenced by the initial prefill plasma density. A significant result of these experiments was the demonstration of control of the injector current level through the plasma flow. The tests also provided a confirmation of virtual cathode effects. Extraction of relatively parallel beams through a 5 cm slit (compared to a 0.7 cm gap) was accomplished. By introducing a curvature into the magnetic field lines near the gap, the beam was brought to a line focus 7 cm downstream.

An injector used recently on the Pulselac C accelerator (Lockner, et al., 1981) is shown in Figure 31. In order to transport plasma to the anode, there can be no magnetic field upstream from the gap. A conducting field excluder was used in conjunction with a downstream pulsed coil set to produce the field pattern shown. Plasma was injected through slots in the excluder. Experiments to date have been conducted with Ne^+ and N^+ ions. Beams of over 5 kA with $j_i \sim 20 \text{ A/cm}^2$ have been produced at 0.1 MV for 0.5 μs pulses. A good understanding of trapped electron effects on net ion flow has been achieved. The area of greatest present interest is the optics of beam extraction. Parallelism from this injector is presently in the range ± 3 degrees.

D. Post-Acceleration and Transport

A number of experiments have been performed on the post-acceleration of beams in the gap of Figure 20. The first, shown in Figure 32, was performed on a system with an injector

Figure 31. Detail of Pulselac-C injector gap.

Figure 32. Post-acceleration experiment. A: Power feed for injector gap magnetic field coil. B: Vacuum insulator. C: Oil housing for injector gap voltage from Marx generator 1. D: Voltage feed for injector gap. E: Flashboard plasma source for protons. F: Injector gap magnetic coil. G: Injector gap vacuum housing. H: Injector gap. I: Extracted ion beam. J: Inductively-isolated voltage feed for post acceleration gap from Marx generator 2. K: Post-acceleration gap.

and one inductively isolated gap (Humphries, et al., 1980a). Gap voltages were supplied by two independent Marx generators. Protons and C^+ were studied. A current of 3 kA of C^+ was supplied to the second gap, which operated at 200 kV in a $0.3 \mu s$ pulse. Measurements of beam energy confirmed straightforward acceleration of the full current in the two gaps. The curvature of the applied magnetic field could be varied in the second gap. Electrostatic focusing was demonstrated with a focal length determined by the curvature. Neutralized transport was investigated by measurements of current density at various locations. By comparison with un-neutralized beam predictions (see Section III,A,1), the positive and negative space charge was found to be balanced to within 0.2 percent. A second experiment (Humphries, et al., 1980b) used an injector gap with four magnetically insulated post-acceleration gaps driven by stages of a single Marx generator. Power for the gap coils and plasma guns was supplied through air core isolation inductors. Beams of C^+ were produced with current exceeding 3 kA in a $0.5 \mu s$ pulse at a peak energy of 0.7 MeV. After passing through the 5 gaps, the beam divergence was found to be less than 0.7 degrees.

The Pulselac C accelerator, presently being tested, is an inductive linear accelerator for neutralized beam propagation studies. It is also a test bed for the development of high current accelerator technology. A diagram of the injector, one post-acceleration stage and a transport chamber with toroidal field lens is shown in Figure 33.

Figure 33. Pulselac-C linear inductive ion accelerator with detail of cavity. System shown with one post-acceleration gap. A: Injector voltage supply oil housing. B: 120 kV pulse-forming-network. C: Isolation transformer for plasma gun power and controls. D: Gas injection plasma gun with plasma expansion chamber. E: Injector gap. F: Post-acceleration gap. G: Power feeds for inductive cavity. H: Ferrite isolation cores. I: Toroidal field lens and propagation chamber.

A detailed view of the cavity is given to illustrate the principle of the inductive linear accelerator. Power entering the cavity can take two paths, through the beam load or back through the cavity structure as shown. Normally, the latter path is almost a short circuit, but the addition of fast response toroidal ferromagnetic cores increases the inductance of the leakage path so that it is almost an open circuit. Power can be coupled to the beam to raise it to high energy after passing through many gaps, yet no part of the system is at a higher voltage than that of a single gap. The accelerator is a transformer, with the beam acting as the secondary and the individual gaps acting as multiple primaries. The properties of the cores present no limitations to the current that can be driven in the load, if they have a flux swing, ΔB_s , that satisfies the condition, $\Delta B_s A_c = V_o \Delta t_p$, where A_c is the core cross sectional area, V_o is the gap voltage, and Δt_p is the voltage pulselength.

The Pulselac C cavity uses ferrite isolation cores, and can sustain a 250 kV pulse for 60 ns. A portion of the injector pulse is post-accelerated. Power is supplied by a 12.5 Ω , 250 kV water filled Blumlein line. This pulse modulator is command fired by a trigatron triggered by a 40 kV hydrogen thyatron which is controlled by digital delay circuits. System jitter is less than 5 ns. Power is transferred to the cavity through standard polyethylene high voltage cables (400 kV breakdown). The charging cycle of the Blumlein line makes a prepulse current that automatically biases the ferrite cores to reverse saturation.

The inductive stage has accelerated 3 kA N^+ beams at a 200 kV voltage level. The divergence of the post-accelerated beam was less than 2 degrees. The beam propagated ballistically in a 1 meter, field free drift region. Future plans call for the addition of 5 stages to Pulselac C which will permit the study of longitudinal dynamics under high beam loading conditions. Unequal currents can be supplied to the upstream and downstream coil sets in the gaps to change the field curvature externally. This will allow tuning of the beam transport system. A toroidal field lens with 85 percent transparency has been tested in vacuum to 13 kG.

E. Inertial Fusion Accelerators

To conclude this section, we will discuss parameters of inertial fusion accelerators consistent with the theory and concepts that have been introduced. Design studies have been initiated on Pulselac D (Humphries, et al., 1981c; Humphries and Lockner, 1981), a demonstration system which could provide focused power density exceeding 5 TW/cm². A scale drawing of the accelerator and one of the cavities with ferrite cores is shown in Figure 34. Design parameters are listed in Table I. The system can achieve high power through a straightforward extrapolation of the parameters of Pulselac C. Current density is within a factor of three of that already achieved and power conditioning could be performed at the 250 kV level, with MV gap potentials obtained by radial stacking (Leiss, 1979). The high current accelerator approach has a number of virtues with

Figure 34. Pulselac-D Fusion Test System. Scale drawings of accelerator and detail of one cavity.

respect to inertial confinement fusion, including high efficiency, clearly defined extrapolation paths, and modularity. Recent advances in the development of glass-like ferromagnetic material will make it possible to produce isolation cores with flux swing four times that of ferrites (VanDevender, 1981). The cross sectional core area shown in Figure 34 could be reduced accordingly.

TABLE I

Parameters of PULSELAC D, Inertial Fusion Test Accelerator.

Current (before bunching)	100 kA
Voltage per stage	1 MV
Number of stages	10
Output energy	10 MeV
Pulselength (before bunching)	60 ns
Bunching factor	1-5
Ion species	He ⁺
Drift tube diameter	60 cm
Source flux	70 A/cm ²
Source divergence	1.5°
Source voltage	0.2 MV
Exit divergence	0.3°
Final focusing lens focal length	0.5 m
Focusing lens maximum field	20 kG
Focal spot diameter	0.5 cm
Focused beam power	5-25 x 10 ¹² W/cm ²
Accelerator length	8 m
Bunching length	3.5 m

VI. CONCLUSIONS

Emphasis in this paper has been on theories of high current ion beam acceleration and transport. Attention has been devoted mainly to the process of vacuum auto-neutralization, in which externally produced electrons are drawn into the ion beam by space charge fields. We will summarize the major results.

A. Electrons from external regions can enter an ion beam and set up distributions which provide almost complete space charge cancellation. This can occur in the absence of collisions and other inelastic processes. To treat the problem, theories must include relaxation pathways for the electrons.

B. Vacuum neutralization has inherent limits. The time response is determined by the rate at which electrons can flow into the beam. The effects of the pressure of the neutralizing electrons can be significant in focused light ion beams.

C. High density ion beams (which cannot propagate without neutralization) can be controlled by confinement of the neutralizing electrons. The electrons can be influenced by magnetic fields too weak to focus directly low velocity ions.

D. Acceleration gaps using radial magnetic field insulation can provide good electron confinement. The magnetic fields introduce no net deflection of paraxial ions. Trapped electrons determine virtual electrode surfaces that make electrostatic focusing by acceleration field possible.

E. Trapped electrons can alleviate space charge effects in acceleration gaps. This is particularly important in injector gaps. Theoretical studies and experiments have demonstrated ion flow enhanced more than an order of magnitude above the Child limit.

F. High current accelerators, with neutralized or bare beams, must provide beam transport in the presence of non-linear forces acting in the transverse direction. Beams can be propagated in this circumstance with little emittance growth if there are many focusing lenses per transverse oscillation (small phase advance per focusing element).

G. Longitudinal confinement and stability is one of the most fundamental problems to be solved for high current linear induction ion accelerators.

Experimental work has been performed with system parameters motivated mainly by inertial fusion applications. Major results are as follows.

A. Progress has been made toward the development of practical intense ion beam injectors. The present generation uses active anode plasma sources to achieve high beam purity, control of injector impedance, and high repetition-rate capability.

B. Pulsed electron sources have been studied. The surface plasma sources developed provide an alternative to thermionic or field-emission cathodes.

C. The behavior of intense ion beams in post-acceleration gaps has been shown to be controllable and scalable. Neutralized transport methods are compatible with practical accelerator technology, such as the ferromagnetic core linear induction accelerator.

The development of high-current, multi-stage accelerators for use as inertial fusion drivers has progressed steadily and has been accompanied by an increased understanding of neutralization mechanisms. At present, there are no technical factors preventing the construction of high flux fusion facilities. New knowledge of trapped electron effects in magnetically insulated gaps suggests the possibility of high power generators operating at lower field stress for long pulses. Compact accelerators operating in the range 100-1000 A for millisecond pulses are feasible. At these pulselengths, temporal limits on vacuum neutralization are not significant and the theories presented are directly applicable. High average power ion accelerators have numerous applications to industrial processes and magnetic confinement fusion. At present, there is a good understanding of intense beam physics in vacuum acceleration gaps. The main challenge for future development of long pulse injectors is the plasma source.

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Advanced Power Systems Division
Electric Power Research Institute
3412 Hillview Ave.
Palo Alto, CA 94304

Prof. G. Bekefi
Research Laboratory of Electronics
Massachusetts Institute of
Technology
Cambridge, MA 02139

James Benford
Pulsed Power Systems Division
Physics International
2700 Merced Street
San Leandro, CA 94577

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Stanford, CA 94305

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Cyclotron Laboratory
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East Lansing, MI 48824

W. Cooper
Lawrence Berkeley National
Laboratory
Berkeley, CA 94720

Francis Cole
Fermi National Accelerator
Laboratory
P. O. Box 500
Batavia, IL 60190

T. Coffee
Naval Research Laboratory
Code 7700
4555 Overlook Ave., SW
Washington, DC 20375

Gerald Cooperstein
Naval Research Laboratory
4555 Overlook Ave., SW
Washington, DC 20375

Wesley O. Doggett
Department of Physics
North Carolina State University
Box 5342
Raleigh, NC 27650

Adam Drobot
Science Applications, Inc.
8400 West Park Ave.
McLean, VA 22015

A. Dangor
Physics Department
Imperial College
London SW7
United Kingdom

Dr. R. Dei Cas
Commissariat a L'Energie Atomique
Centre D'Etudes de Bruyeres-le-
Chatel
B. P. No. 561
92542 - Montrouge - Cedex France

A. Faltens
Lawrence Berkeley National
Laboratory
1 Cyclotron Road
Berkeley, CA 94720

Moshe Friedman
Naval Research Laboratory
Code 7700.1
4555 Overlook Ave., SW
Washington, DC 20375

Terry F. Godlove
Department of Energy
Mail Code C404(LF)
Washington, DC 20545

Ryszard Gajewski
U. S. Department of Energy
Division of Advanced Energy
Projects
Mail Stop J309
Washington, DC 20545

Dr. Ernst Granneman
FOM Institute for Atomic and
Molecular Physics
Amsterdam
The Netherlands

Zaven Guiragossian
TRW Systems and Energy
One Space Park
R1/1078
Redondo Beach, CA 90278

Brendan Godfrey
Mission Research Corporation
1400 San Mateo Blvd. SE
Albuquerque, NM 87110

Lt. Col. Anthony Hyder
Air Force Office of Scientific
Research - Nuclear Physics
Bolling Air Force Base, DC 20332

Ingo Hoffman
Max-Planck-Institut für
Plasmaphysik
8046 Garching bei München
Federal Republic of Germany

D. A. Hammer
Laboratory of Plasma Studies
308 Upson Hall
Cornell University
Ithaca, New York 14850

Dr. K. Ikuta
Research Information Center
Institute of Plasma Physics
Nagoya University
Nagoya, Japan

Hiroshi Ishizuka
(Associate Professor)
Institute of Physics
University of Tsukuba
Ibaraki, 300-31 Japan

Robert A. Jameson
Los Alamos National Laboratory
P. O. Box 1663, MS 820
Los Alamos, NM 87544

Werner Joho
Schweizerisches Institut für
Nuklearforschung
5234 Villigen
Switzerland

P. K. John
Department of Physics, Faculty
of Science
Physical Sciences Building
University of Western Ontario
London, Ontario
Canada

W. B. Kunkel
Lawrence Berkeley National
Laboratory
Berkeley, CA 94720

Prof. Sunao Kawasaki
Department of Physics
Kanazawa University
1-1 Marunouchi
Kanazawa, Ishikawa, 920
Japan

Bernard Kulke
FXR Project
Lawrence Livermore National
Laboratory
P. O. Box 808, MS L-153
Livermore, CA 94550

Denis Keefe
Lawrence Berkeley National
Laboratory
Berkeley, CA 94720

N. Marshall King
SRC Rutherford Laboratory
Chilton, Didcot
Oxon. OX 11 0QX
United Kingdom

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Institut für Angewandte Physik
Johann Wolfgang Goethe -
Universität
Robert-Mayer-Strasse 2-4
D-6000 Frankfurt am Main 1
West Germany

Professor Dr. Wolfgang Klöse
Mitglied des Vorstandes
Kernforschungszentrum Karlsruhe
GMBH
D-7500 Karlsruhe
Germany

R. Kusse
Laboratory of Plasma Studies
Upson Hall
Cornell University
Ithaca, NY 14853

J. D. Lawson
Science Research Council
Rutherford Laboratory
Chilton, Didcot, Berkshire
OX 11 0QX
United Kingdom

Don S. Lemons
Intense Particle Beam Theory
Group
Los Alamos National Laboratory
Los Alamos, NM 87544

Alfred Mondelli
Science Applications, Inc.
8400 West Park Ave.
McLean, VA 22015

A. W. Maschke
Brookhaven National Laboratory
Upton, NY 11973

R. V. Pyle
Lawrence Berkeley National
Laboratory
Berkeley, CA 94720

Dr. Andreas Pritzker
SIN
Schweizerisches Institut für
Nuklearforschung
5234 Villigen
Switzerland

Dr. C. Patou
Laboratoire D'Environnement
Neutronique et Photonique
Centre de Valduc
Dijon, France

D. Prono
Lawrence Livermore National
Laboratory
P. O. Box 808, MS L-637
Livermore, CA 94550

Steven Rosenblum
Lawrence Berkeley National
Laboratory
Bldg. 47
Berkeley, CA 94720

Scott Robertson
Physics Department
University of California
Irvine, CA 92717

Martin Reiser
Department of Electronics
Engineering
University of Maryland
College Park, MD 20742

Norman Rostoker
Physics Department
University of California
Irvine, CA 92717

W. Schmidt
Kernforschungszentrum
P. O. Box 3640
D-7500 Karlsruhe
West Germany

John Siambis
Naval Research Laboratory
Plasma Physics Division
Washington, DC 20375

Jeremy Stein
Pulsar Associates
5000 Copper NE
Albuquerque, NM 87108

Waldemar Scharf
Warsaw Technical University
Institute of Radioelectronics
Nowowiejska 15/19
00-665 Warsaw
Poland

R. N. Sudan
Laboratory of Plasma Studies
Upson Hall
Cornell University
Ithaca, NY 14850

Michael J. Shaw
Rutherford Laboratories
Chilton, Didcot, Oxon OX 11 0QX
United Kingdom

A. C. Smith, Jr.
Room 2518
Pacific Gas and Electric Company
77 Beal Street
San Francisco, CA 94106

A. Toepfer
Physics International Company
2700 Merced Street
San Leandro, CA 94577

W. L. Talbert
Los Alamos National Laboratory
P. O. Box 1663, MS 811
Los Alamos, NM 87544

F. Winterberg
Desert Research Institute
University of Nevada
Reno, NV 89506

H. Wollnik
II. Physikalisches Institut der
Universität
D-63
Giessen, West Germany

Mark Wilson
Bldg. 245, B102
U. S. Department of Commerce
National Bureau of Standards
Washington, DC 20243

Paul Wheeler
Lawrence Livermore National
Laboratory
A Division
Livermore, CA 94550

Kiyoshi Yatsui
Laboratory of Beam-Fusion
Technology
Tech. Univ. of Nagaoka
Nagaoka, Niigata 949-54
Japan

K. Zieher
Kernforschungszentrum
P. O. Box 3640
D-7500 Karlsruhe
West Germany

Dr. Chiyoe Yamanaka
Institute of Laser Engineering
Osaka University
Yamada-kami, Suita
Osaka, 565, Japan

4000	A. Narath
4200	G. Yonas
4241	J. R. Freeman
4241	C. L. Olson
4241	J. W. Poukey
4241	J. P. Quintenz
4241	J. A. Swegle
4242	K. W. Bieg
4242	P. L. Dreike
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4244	G. S. Mills
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4244	S. E. Rosenthal

4247 S. A. Slutz
4250 T. H. Martin
4252 J. P. VanDevender
4253 K. R. Prestwich
4253 J. J. Ramirez
4255 R. B. Miller
4255 S. Humphries (20)
4255 T. R. Lockner (10)
4255 M. G. Mazarakis
4443 J. Sprung
8265 P. Dean
3141 L. J. Erickson (5)
3151 W. L. Garner (3)
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