Handwritten: R107-452

MAGNETICALLY PERTURBED PLASMA CONFINEMENT
IN A TOROIDAL OCTUPOLE INCLUDING OBSTACLE INFLUENCE

MASTER

J. R. Drake and J. R. Greenwood

University of Wisconsin, Madison, Wisconsin 53706

Permanent-magnet dipole perturbations were positioned outside the external conducting wall of a pulsed, shearless toroidal octupole plasma containment device. Floating potential profiles and losses were studied for two perturbation configurations; field lines were either closed within the plasma containment zone or were drawn outside the containment zone. Profiles were not affected by the former configuration but were altered when the field lines were drawn outside the containment zone. The extent of the profile change was enhanced when the azimuthal $\bar{E} \times \bar{B}$ drifts were suppressed by the presence of obstacle-like internal ring supports. With the $\bar{E} \times \bar{B}$ drifts suppressed, the incidence of repetitive encounters was increased between transit particles following field lines and the perturbation field. The net result was particle trajectory displacement and altered V_f profiles.

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I. INTRODUCTION

Various experiments have demonstrated that magnetic field errors can cause enhanced plasma transport from the containment zone in toroidal multipole containment devices with purely-poloidal magnetic fields. Distinctions must be made between magnetic perturbations that leave field lines contained, and for some cases even closed, within the plasma containment zone and perturbations that draw lines from the containment zone, along a spiral course, into either the internal ring surfaces or the external wall surfaces. The usual magnetic perturbations found in pulsed, conducting-wall octupole experiments are examples of the contained-line case. For this case, plasma transport can be the result of gradient drifts due to the magnetic perturbation,^{1,2} or drifts around magnetic islands.³ On the other hand, uncontained field lines are normally associated with containment fields and magnetic perturbations that are constant in time. Flow along the spiral course of an uncontained field line can enhance plasma transport either directly or secondarily through electric field drifts due to plasma potential asymmetries generated by the magnetic perturbation.⁴⁻⁷ Uncontained lines can be created in machines with pulsed fields and conducting boundaries by applying permanent magnets outside the conducting boundary.

In experiments previously reported, the internal rings of the multipole devices were mechanically supported by obstacle-like prods or wires which passed through the containment zone. The experiments to be

described here were performed on a pulsed octupole with a conducting external wall. During the containment field pulse, the ring supports and all other connections and obstacles could be removed.⁸ Permanent-magnet, dipole perturbations were applied outside the external conducting wall of the toroid while operating in the purely-poloidal field mode and the plasma was scanned for floating potential structure.

The boundary conditions for these magnetic perturbations did not match the boundary conditions on the pulsed main field at the conducting wall. The field topology then was not an example of that which occurs with inherent field errors from distortions in the conducting wall. However, in one phase of the experiment, two dipole perturbations were positioned and oriented so that the perturbed field lines were both closed and contained within the plasma containment zone. This closed-line configuration was similar to topologies which result from certain types of frequently encountered wall distortions such as access holes positioned with the proper symmetry. The floating potential profiles for this configuration were identical to those observed when no perturbations were applied, although axisymmetry was destroyed by these two, compensating perturbations. This suggests that the effects of magnetic field errors can be minimized by arranging construction-related distortions in the conducting wall so that field lines are left closed and contained within the plasma containment zone wherever possible.

In addition, an uncontained field line topology was generated by applying one permanent-magnetic dipole perturbation. When the internal

ring supports were present in the plasma containment zone, the floating potential contours exhibited two cells of opposite vorticity, a pattern similar to that seen in other dipole-perturbed, supported multipole devices.^{5,9} However, when the supports were removed, with the magnetic perturbation still in place, the two-cell contour pattern was not seen and the contours exhibited only minor azimuthal asymmetries.

An examination of the observed azimuthal drifts for single particles, including $\bar{E} \times \bar{B}$ drifts inferred from the spatial dependence of the floating potential and ∇B drifts, provided an explanation of the observed difference between the cases with and without ring supports in the plasma containment zone. When no ring supports were present, the nearly axisymmetric azimuthal $\bar{E} \times \bar{B}$ drift was such that only those electrons in the high energy part of the electron distribution were both untrapped (circulating along closed poloidal field lines around the minor axis of the toroid) and azimuthally drifting with a near-zero net angular velocity. Untrapped electrons were capable of parallel flow along uncontained field lines⁷ and therefore would be expected to give rise to significant transport if their azimuthal drift velocity was small. This transport would be larger than predicted transport due to magnetic-perturbation gradient drifts by a factor of the order of the ratio of the perturbation size to the gyroradius.² When ring supports were present in the plasma, the azimuthal $\bar{E} \times \bar{B}$ drift velocity was reduced. As a result there was an increase in the estimated fraction of the electron population that was both untrapped and azimuthally drifting with a near-

zero net angular velocity. Therefore the cumulative effects of the magnetic perturbation were increased when the ring supports were present as manifested in observed asymmetries in the floating potential profile.

II. APPARATUS AND PLASMA PARAMETERS

In the Wisconsin octupole, the magnetic flux function ψ is defined such that the poloidal flux is equal to $2\pi\psi$. An orthogonal coordinate system χ , ψ and θ is used where θ is the toroidal angle and χ is the magnetic scalar potential. The magnetic field is given by $\vec{B} = \vec{\nabla}\chi = \vec{\nabla}\psi \times \vec{\nabla}\theta$ and the internal ring current is in the direction of decreasing θ . The flux function notation for the separatrix is ψ_s . The last stable flux surface outside ψ_s is indicated as ψ_c ($\oint B^{-2} d\chi$ is minimum), and the flux surface where the field line length $L = \oint B^{-1} d\chi$ is minimum is indicated as ψ_L .

The field pulse has a sinusoidal time variation with a half-sine wave period of 43 msec. However in these experiments the condensor bank was shorted (crowbarred) 25 msec after the start of a field pulse and the magnetic field then decayed with a decay time of about 70 msec. A flux plot 25 msec after the start of the field pulse is shown in Fig. 1. At this time the separation between a ring surface and ψ_c was about equal to thirty 20-eV ion gyroradii. The maximum field strengths at the internal ring surfaces were about 5.3 kG for the two rings which have the smaller major radius and 2.5 kG for the two rings which have the larger major radius.

The plasma was the afterglow of a microwave-generated plasma. Starting with a background pressure of 2×10^{-7} Torr, hydrogen was injected to establish a neutral H_2 base density of about 10^{11} cm^{-3} before pulsing the magnetic field and activating the microwaves. After partial

preionization with a 100-W, continuous microwave source, a 144- μ sec pulse of 100-kW microwaves was activated for heating¹⁰ and additional ionization of the hydrogen. The parameters of the afterglow plasma were $n = 5 \times 10^9 \text{ cm}^{-3}$, $kT_e = 2\text{-}10 \text{ eV}$, and $kT_i = 10\text{-}40 \text{ eV}$. The mean free paths between electron-ion and electron-neutral collisions were both about one-hundred times greater than the major radius of the toroid 3 msec after the microwave heating pulse.

The presence of supports affected the floating potential profiles; the ratio of the total ring support (obstacle) surface area to containment volume was $8.3 \times 10^{-5} \text{ cm}^{-1}$. This gave a 20-msec inverse loss rate due to sound speed flow to all the support obstacles.

Langmuir probes were used to measure the spatial dependence of the plasma parameters. The electron temperature profile was quite flat as a function of ψ . Therefore electric fields were inferred directly from floating potential profiles.

When a magnetic perturbation is applied, the perturbed field line has a net perpendicular deviation from the flux surface of the unperturbed field line for each circulation around the minor axis. The deviation is conveniently measured in the units of the unperturbed flux function. If the perpendicular component of the field error is δB_ψ , the net deviation is given by $\Delta\psi = \oint R \delta B_\psi d\ell$ where R is the major radius at the field line element $d\ell$. Therefore the deviation of a field line perpendicular to a ψ -surface is given by the perturbation magnetic flux that penetrates the ψ -surface per azimuthal radian.

In this paper, two magnetic dipole perturbations are referred to as compensating if field lines remain closed within the containment zone. Each dipole perturbation causes a deviation in the field lines, $\Delta\psi_k$, and certainly $\Delta\psi = \sum \Delta\psi_k = 0$ is a necessary criterion for perturbations to be compensating. Examples of two compensating perturbations are the two identical dipoles shown located at Q and P in Fig. 1 which are coaxial and symmetrically located above and below the horizontal symmetry plane of the toroid.

These dipole perturbations are described as compensating because, in the collisionless case, there is no net plasma transport derived from parallel flow along the perturbed field line. However for transport derived from magnetic-perturbation gradient drifts, the possibility of a net cancellation of the drifts depends on the geometry of the system. If $\bar{B}_0$ and δB are the unperturbed field and perturbation, one contribution to the ψ component of the gradient drift that is first order in δB is proportional to $B_0^{-2} \nabla_{\theta} \delta B_x$. The integral of this quantity along an unperturbed field line is small if the axis of the dipole is perpendicular to $\bar{B}_0$ and if $\bar{B}_0$ is nearly constant (or has the proper symmetry) in the region where δB_x is large because of the odd symmetry of $\delta B_x(x)$. Therefore, if the parallel velocity of the particle is nearly constant and much greater than the azimuthal drift velocity in the vicinity of the perturbation, this component of the perturbation gradient drift can be described as self-compensating.

The general form for the ψ -component of the first order drift

for the systems discussed here is proportional to $B_0^{-2} \delta B \partial B_0 / \partial \chi$.

Transport derived from this drift can be compensated by felicitous placement of dipole perturbations in pairs as is the case with parallel flow transport.

III. ANALYSIS OF AZIMUTHAL DRIFTS

The drift velocity is given by

$$\bar{v}_d = \frac{m}{q} \left(\frac{v_\perp^2}{2} + v_\parallel^2 \right) \frac{\bar{B} \times \nabla \bar{B}}{B^3} - \frac{\nabla \phi \times \bar{B}}{B^2}, \quad (1)$$

where $v_\perp$ and $v_\parallel$ are the components of the particle velocity perpendicular and parallel to B , m/q is the mass to charge ratio, and ϕ is the plasma space potential. The time-averaged angular velocity ω_d due to the drift in the azimuthal direction in a toroidal multipole is given by

$$\omega_d = \frac{1}{\tau_b} \int_0^{\tau_b} \frac{v_d \cdot \hat{\theta}}{R} dt, \quad (2)$$

where $\tau_b = \oint dl/v_\parallel$ is the bounce period for a particle trapped in one of the mirrors that exist in a multipole or the transit period for an untrapped particle circulating around the minor axis of the toroid. From Eqs. (1) and (2), for the axisymmetric case with a purely-poloidal B and ϕ constant along a field line, ω_d is given by

$$\omega_d = \frac{\partial \phi}{\partial \psi} + \frac{1}{\tau_b q} \oint \left(\frac{mv_\parallel}{B} + \frac{\mu}{v_\parallel} \right) \frac{\partial B}{\partial \psi} dl, \quad (3)$$

where $\mu = (\frac{1}{2}mv_\perp^2)/B$. The $\bar{E} \times \bar{B}$ drift contribution given by $\partial \phi / \partial \psi$ is in the same direction for all charged particles. But the direction of ω_d , the $\bar{v}B$ -drift contribution to ω_d given by the second term on the right-hand side of Eq. (3), depends on whether the integral in the term is dominated by contributions from the unfavorable curvature region ($\partial B / \partial \psi < 0$) or the favorable curvature region ($\partial B / \partial \psi > 0$), as well as

the charge of the particle.¹¹

For a given ψ and $K = \frac{1}{2}mv^2$, the range of ω_d' with varying μ is bounded. The extrema of ω_d' are summarized in Table I. Trapped particles with turning points ($v_{||} = 0$) at points along the field lines where $B^{-1}\partial B/\partial\psi$ is an extremum have ω_d' near the extrema values of ω_d' . A circulating particle with $\mu = 0$ has $\omega_d' = - (2K/q)(L^{-1}\partial L/\partial\psi)$ which is zero at ψ_L .

In the range $\psi_s < \psi < \psi_c$, the minimum of $B^{-1}\partial B/\partial\psi$ as a function of χ for a constant ψ is nearly independent of ψ . However the value of the maximum decreases with increasing ψ and is roughly proportional to $(\psi - \psi_s)^{-1}$. Combining the range of ω_d' expressed in Table I with the azimuthal $\bar{E} \times \bar{B}$ drift contribution to ω_d yields a range of possible values of ω_d for a given ψ and K .

In this analysis we assume that enhanced transport due to the effects of magnetic perturbations is initiated by displacements of the particle guiding-center. A critical azimuthal angular velocity $\omega^*(K, \mu, \psi, \delta B)$ is defined such that if we have $|\omega_d| > \omega^*$ then the ψ displacement of the particle rapidly drifting past the perturbation is less than ρ_0 and can be considered negligible. The critical displacement used here is an ion gyroradius calculated using the ion thermal speed and a field strength equal to $L/(\oint dL/B)$ evaluated at ψ_L . A criterion based on ρ_0 is somewhat arbitrary but is useful in depicting the dependence of ψ displacement on ω_d . Then $W_0/(R\omega^*) = \rho_0/V_\psi$ is the maximum time a particle can take to azimuthally drift past the pertur-

bation and still not experience a displacement greater than ρ_0 . Here we have used W_0 , the azimuthal width of the perturbation, and $V_\psi(K, \mu, \psi, \delta B)$, the velocity of the guiding center in the ψ direction time-averaged over τ_b .

The contribution to V_ψ from parallel flow along the spiral course of a field line can be distinguished from that due to gradient drifts. For the perturbation systems discussed here the dominant magnetic-perturbation gradient drift is proportional to $B_0^{-2} \delta B \partial B_0 / \partial \chi$. This drift is compensated for the dipole perturbations shown in Fig. 1 where $\Delta\psi = \Delta\psi_Q + \Delta\psi_P = 0$ unless the particle is trapped with trajectory limitations that prevent the particle from experiencing the compensating effects. In the discussion that follows we use the notation $\Sigma \Delta\psi_k$ to indicate that particle trapping must be considered in evaluating transport due to perturbation gradient drifts.

Because V_ψ is the sum of contributions resulting from parallel flow and from magnetic-perturbation gradient drifts, we can consider $\omega^* = (V_\psi / \rho_0)(W_0 / R)$ as having two corresponding components, $\omega_{\parallel}^*$ and ω_V^* . Detailed calculations of the particle trajectories are necessary to calculate V_ψ and ω^* exactly. The following approximations are used here,

$$\omega_{\parallel}^* \approx \frac{W_0}{\tau_b \rho_0} \left[\frac{\Delta\psi}{R^2 B} \right] \quad (4)$$

and

$$\omega_V^* \approx \frac{W_0}{\tau_b \rho_0} \left[\frac{\partial B}{\partial \chi} \frac{\rho^*}{R^2 B} \Sigma \Delta\psi_k \right], \quad (5)$$

where $\rho^* = (\frac{1}{2}mv_{\perp}^2/v_{\parallel} + mv_{\parallel})/(eB)$ and the brackets indicate that the enclosed parameters are evaluated using typical values in the vicinity of the perturbation. If $v_{\parallel}$ is near zero at the perturbation, Eq. (5) reduces to

$$\omega_{\nabla}^* \approx \frac{W_{\theta}}{2W_{\chi}\rho_0} \left[\frac{\partial B}{\partial \chi} \frac{\rho_{\perp} v_{\perp}}{R^2 B} \sum \Delta \psi_k \right]$$

where W_{χ} is the perturbation size in the χ direction and $\rho_{\perp} = (mv_{\perp})/(qB)$. The effect of the magnetic perturbation is enhanced for trapped particles with turning points coincident with the perturbation.

For a single dipole perturbation (nonzero $\Delta\psi$) with $\rho \ll W \leq (\partial B/\partial \chi)^{-1}$ we have $\omega_{\nabla}^* \gg \omega_{\nabla}^*$. By our definition of ω^* , we can have ψ displacement greater than ρ_0 as the particle azimuthally drifts past the perturbation if we have $|\omega_d| = |\partial\phi/\partial\psi + \omega_d'| < \omega_{\nabla}^*$. This inequality is represented for electrons in Fig. 2 for $\psi = \psi_L$. All ω_d fall between the extrema of ω_d' , and the intersection of this area with the area between $\pm \omega_{\nabla}^*$ represents electrons affected by the perturbation. As can be seen in Table I, the curves of the ω_d extrema for ions can be obtained by interchanging maximum (minimum) values of ω_d' for electrons with the negative of the minimum (maximum) values of ω_d' for ions.

The magnitude of $\Delta\psi/(R^2 B)$ used in evaluating ω_{∇}^* in Fig. 2 was such that the spiral course of the perturbed field line encircled the minor axis of the toroid 100 times between ψ_s and ψ_c . This is considered a 1% field perturbation. Other values used were $kT_e = 10$ eV, $W_{\theta}/\rho_0 = 20$ and $\partial\phi/\partial\psi = - (kT_e/e)/(\psi_c - \psi_s)$. As is discussed in Sec. IV, this value

for $\partial\phi/\partial\psi$ is representative of that observed near ψ_L for the unperturbed, support-free case where $\partial\phi/\partial\psi$ is fairly axisymmetric.

According to Fig. 2, the electrons with drift trajectories susceptible to a ψ displacement greater than ρ_0 are those with energies greater than K_e^* . Any circumstance that results in a decrease in $|\partial\phi/\partial\psi|$ in turn lowers K_e^* thus increasing the energy range corresponding to electrons susceptible to large trajectory displacements. On the other hand, selecting a cutoff for significant ψ displacement at some value greater than ρ_0 , decreasing W_0/ρ_0 , or decreasing $\Delta\psi/(R^2B)$ would decrease the number of electrons susceptible to significant ψ displacement. In Sec. IV the effect of a decreased $|\partial\phi/\partial\psi|$ is discussed in detail because the presence of ring supports decreased $|\partial\phi/\partial\psi|$ and the magnetic perturbation induced disturbance to the plasma was enhanced.

For the cases considered here, $|\delta B|$ is negligible for flux surfaces inside a surface about midway between ψ_L and ψ_S . In the region where $|\delta B|$ is significant, we have $|L^{-1}\partial L/\partial\psi|$ less than both $|B^{-1}\partial B/\partial\psi|_{\max}$ and $|B^{-1}\partial B/\partial\psi|_{\min}$. Therefore the plots presented in Fig. 2 for $\psi = \psi_L$ are representative of the entire ψ range of interest.

Curves indicating $\pm \omega_H^*$ for ions are also shown in Fig. 2. Interchange of the ω_d extrema curves leads to an analogous minimum ion energy limit K_i^* . For a given K , the region between the extrema of $\pm \omega_H^*$ is very small for ions. The anticipated ion transport derived from ψ displacement of single particle trajectories is therefore much less than electron transport for the parameters discussed here.

If a second dipole is added with the same $|\Delta\psi_k|$ so that $\Delta\psi = \Sigma\Delta\psi_k = 0$, then $\omega_{\parallel}$ is zero and $\pm\omega_{\nabla}^*$ would determine K_{θ}^* and $K_{\parallel}^*$. But $|\omega_{\nabla}^*|$ is much less than $|\omega_{\parallel}^*|$ for a single dipole. Therefore the low energy limit of the energy range of those particles susceptible to transport is high thus excluding most particles. Furthermore, only particles trapped in the mirror where the perturbation is localized experience non-compensating drifts in the ψ direction.

The anticipated effect of increased collisionality is a decrease in the drift trajectory displacement derived from parallel flow along a nonzero- $\Delta\psi$ perturbed field line.⁷ About 80%-90% of the particles between ψ_L and ψ_C are trapped. Therefore small angle scattering events limit the time a circulating electron exists in an untrapped state to an effective collision period τ_c which is less than the collision period for 90° scattering. As the plasma cools, the increase in collisionality decreases the transport of plasma derived from parallel flow.

It is also necessary to consider whether the effects of the perturbation can alter ω_d sufficiently to invalidate the above considerations which included only an examination of the unperturbed ω_d . The displacement of the particle trajectory in the θ direction, Δx_θ , must be considered as well as the displacement in the ψ direction, Δx_ψ . If the effects of the perturbation gradient drifts alone are considered, then we have $\Delta x_\theta \approx \Delta x_\psi < \rho_0$ and ω_d is not significantly altered by the perturbation. The drifting particle cannot be stalled at

the azimuth of the perturbation by perturbation gradient drifts alone. However the asymmetric potential, $\delta\phi$, that develops when the perturbation is applied and the associated $\bar{E} \times \bar{B}$ drifts can be significant. These potential asymmetries cause drifts in both the ψ and θ directions; the former affects transport and the second alters ω_d . The observed $\delta\phi$ and the associated drifts are discussed in Sec. IV where experimental cases are discussed.

For many particles, ω_d is sufficiently large for the particles to drift past the perturbed azimuth without experiencing serious ψ displacement. However for the experimental time scales and regions of ψ space considered here, these particles do not drift around the major axis of the toroid at a high enough rate to undergo ordered transport due to repeated encounters with the perturbed azimuth.

IV. EXPERIMENTAL CONCLUSIONS

In the first experiment, the V_f structure generated by a magnetic perturbation with $\Delta\psi$ nonzero (transport dominated by parallel flow) was compared with the case of a perturbation with $\Delta\psi = 0$ (transport solely derived from gradient drifts). A magnetic dipole was first located at position Q, $\Delta\psi = \Delta\psi_Q$, and then position P, $\Delta\psi = \Delta\psi_P$; both these cases have a nonzero $\Delta\psi$. In the second phase, dipoles of equal magnitude were positioned at Q and P such that $\Delta\psi = \Delta\psi_Q + \Delta\psi_P = 0$. Langmuir probe scans of V_f versus θ are shown in Fig. 3. The probe measuring V_f was located in the local mirror region near position Q and midway between ψ_L and ψ_C . The magnetic perturbations were applied at $\theta = 52^\circ$. Ring supports were present and were located at $\theta = 42^\circ$ in the azimuthal segment shown in Fig. 3. At the azimuth coincident with the axis of the dipole perturbation the magnitudes of the perturbations at ψ_C were such that $|\Delta\psi_Q| = |\Delta\psi_P|$ was 1% of $(\psi_C - \psi_S)$. The case represented in Fig. 2 corresponds to a perturbation with this magnitude and size.

As seen in Fig. 3a, the V_f profile was the same whether the single perturbation was located at Q or P. But, as seen in Fig. 3b, the V_f profile observed when compensating perturbations ($\Delta\psi = 0$) were applied was the same as that observed when no perturbations were applied. The case selected to be shown here was particularly striking because the support-generated V_f asymmetry, evident at $\theta = 42^\circ$ in Fig. 3b, was dramatically altered when a nonzero- $\Delta\psi$ perturbation was

applied (Fig. 3a). This asymmetry was restored to the pattern seen when no perturbation was applied by the addition of a second compensating perturbation.

There were four mirrors in series along the field lines in the psi-space region between ψ_s and ψ_c ; one local field maximum was associated with each of the four internal rings. Positions Q and P were separated by two of the maxima in either direction along the field line. The mirror ratio of the local field maximum near an inner (small-R) ring to the field magnitude near Q (or P) for $\psi_L < \psi < \psi_c$ was about 4:1 and the corresponding ratio for an outer (large-R) ring was about 2:1. Therefore about 70% particles in the immediate vicinity of the perturbation were locally trapped in the mirror near Q (or P).

Figure 3a demonstrates that the effects of the magnetic perturbation were not isolated by the mirror throats although 70% of the potentially affected particles were trapped in the local mirrors near Q (or P). Interaction between these trapped particles and the magnetic perturbation did not lead to any observable change in V_f structure. Furthermore, applying a nonzero- $\Delta\psi$ perturbation affected the V_f structure but applying a perturbation with $\Delta\psi = 0$ had no effect. Therefore no observed changes in the V_f structure could be related to the effects of perturbation gradient drifts. The resolution of these scans was about 5% of kT_e/e .

This was consistent with what was outlined in Sec. III where it was shown that, for this perturbation, the magnitude of ω_{∇}^* was such

that only a small fraction of the particle distribution would be susceptible to significant trajectory displacement. However, it is of interest that the estimated magnitude of the local transport due to magnetic-perturbation gradient drifts at the point where δB was largest was of the order of 10% of Bohm diffusion so it is somewhat surprising that no manifestation of this local transport was evident in the V_f profiles. Even the inherent asymmetries generated by ring supports (obstacles) in the plasma containment zone were not sensitive to perturbations with $\Delta\psi = 0$.

An additional measurement of particle transport was made directly using particle collectors. These collectors measure particle fluxes on wall surfaces without significantly perturbing the plasma. The collector was located on the external conducting wall of the toroid near position Q. With the supports present in the containment zone when the nonzero- $\Delta\psi$ perturbation was applied at Q, the particle flux increased to a value a factor of three times greater than that seen with no perturbation. However, when two compensating perturbations were applied at Q and P ($\Delta\psi = 0$), the measured particle flux was the same as that seen with no perturbation.

These results are similar to those reported by Meade et al.,^{5,6} for a supported internal ring device. There the enhanced transport due to magnetic perturbations with $\Delta\psi = 0$ was more than an order of magnitude less than the corresponding enhancement from nonzero- $\Delta\psi$ perturbations. However, in Meade's experiments the perturbation with $\Delta\psi = 0$ was a mag-

netic dipole applied with the dipole axis parallel to the poloidal field lines. A significant difference between the experiments performed by Meade, et al and the experiment described here is the fact that in this experiment $\Delta\psi_Q$ and $\Delta\psi_P$ were isolated from each other by two mirror throats.

To compliment the studies of plasma behavior in the presence of permanent-magnet dipole perturbations which have been discussed so far, a search for perturbation-influenced plasma structure was carried out in the vicinity of a 15-cm diameter access hole in the conducting wall of the toroid. The hole was located in the lower wall with the hole axis perpendicular to the horizontal symmetry plane and midway between the large-R and small-R internal rings (see Fig. 1). The hole could be "filled" by press-fitting a mating, truncated-conical, conducting plug¹² which, when engaged, made good electrical contact with the hole rim. When the plug was not in good electrical contact, $|\delta B|$ was about 5% of $|B_0|$ at ψ_c but $|\delta B|$ was reduced by a factor of 20 when good contact was established. Scans of V_f showed no change when the perturbation was "applied" by breaking contact between the plug and the hole rim.

If toroidal effects were neglected, this wall hole was a perturbation with $\Delta\psi = 0$. There was an apparent similarity between the wall-hole perturbation and the two compensating dipoles with $\Delta\psi = 0$. The field lines were nearly closed for the wall-hole perturbation and perfectly closed for the compensating-dipole perturbation. No perturbation induced structure was observed for either case.

Contour plots of V_f on a ψ versus θ grid are now described that were observed when a nonzero- $\Delta\psi$ magnetic perturbation was applied. In this experiment $|\Delta\psi| = |\Delta\psi_0|$ was about 2% of $(\psi_c - \psi_s)$ at ψ_c and the perturbation was applied at an azimuth coincident with a set of outer (large-R) ring supports. Perturbation effects on the contours with the supports withdrawn from the containment zone are compared to a case with the supports present. A similar supported case is shown in Fig. 3b but with the perturbation 10° from the supports.

Floating potential contours observed 3 msec after the microwave heating and ionizing was completed (28 msec after the start of the containment field pulse) are shown in Fig. 4 for three cases. The cases are as follows: a) supports present and no applied magnetic perturbation, b) supports present with an applied magnetic perturbation, and c) supports withdrawn from the containment zone with an applied magnetic perturbation. The contours in Fig. 4a show the effects of a support obstacle.¹³ In Fig. 4b a two-cell contour is evident and was generated when the magnetic perturbation was applied with the supports present. However only minor differences were observed between the support-free cases with and without a magnetic perturbation. Only the magnetically perturbed support-free case is presented here in Fig. 4c.

When the supports were present in the containment zone (Fig. 4a), the V_f profiles were not axisymmetric. Overall we had $|\partial\phi/\partial\psi| < (kT_e/e)/(\psi_c - \psi_s)$ and $\partial\phi/\partial\psi$ was near zero at the azimuth of the supports for a large ψ interval as inferred from the V_f contour plots. However, when the supports were withdrawn, $\partial\phi/\partial\psi$ was quite axisymmetric

and was of the order of $-(kT_e/e)/(\psi_c - \psi_s)$ for much of the region between the ψ -surface where V_f was maximum and ψ_c . For both cases we had $kT_e \approx 10$ eV at the time shown in Fig. 4 and $\tau_c/\tau_b \approx 50$. The perturbation magnitude was twice that used in the calculations represented in Fig. 2 but the remaining parameters were the same. The supports suppressed the magnitude of $\partial\phi/\partial\psi$. This in turn increased the number of particles in jeopardy of experiencing large ψ displacement through parallel flow while under the influence of the perturbation and the two-cell structure seen in Fig. 4b was the result.

This nonaxisymmetric two-cell structure is different from the nonaxisymmetric structure seen in Fig. 4a when the supports alone were present. Apparently new drift trajectories evolved as a result of the magnetic perturbation-induced ψ displacement of the electron trajectories. In some regions the presence of the magnetic perturbation increased $|\partial\phi/\partial\psi|$ to values greater than those seen when supports alone were present. In these regions this increase in $|\partial\phi/\partial\psi|$ would lead to a decreased population of untrapped electrons with $|\omega_d| < \omega^*$. Therefore it is believed that an equilibrium was established between transport from the containment zone and azimuthal $\mathbf{E} \times \mathbf{B}$ drifts in a fashion similar to that described by Jernigan, et al.⁶

Cooling of kT_e by the support obstacle could increase the collisionality and decrease the losses. However this was not a strong effect at this time.

Floating potential contours are shown in Fig. 5 for the two supported cases (with and without magnetic perturbation) 12.5 msec after the completion of the microwave heating and ionizing pulse to demonstrate the time evolution of the obstacle-generated and the magnetic perturbation-generated potential structure. As can be seen in Fig. 5a, the topology of the contours generated by the obstacle did not change over the 9.5-msec interval between the representations in Fig. 4a and Fig. 5a. However the two-cell structure, generated by the magnetic perturbation, did decay in time and the contour evolved toward that seen in Fig. 5a. The two-cell structure was generated at early times after the heating pulse but was not sustained as the plasma cooled. The increased collisionality decreased the transport derived from parallel flow. At the times shown in Fig. 5, kT_e was of the order of 2 eV and τ_c was about equal to τ_b for electrons. Therefore the two-cell structure was not sustained.

In summary two points are emphasized. First, destroying axisymmetry did not affect the observed V_f profiles when magnetic perturbations were applied in a manner which maintained a closed field line configuration within the containment zone. Second, the presence of ring supports (obstacles) affected the interaction between the plasma and a nonzero- $\Delta\psi$ magnetic perturbation through support induced changes in the potential structure and, therefore, the azimuthal $\mathbf{E} \times \mathbf{B}$ drift. The supports suppressed the magnitude of $\partial\phi/\partial\psi$ which in turn increased

the population of untrapped electrons which also had near-zero azimuthal drifts. Only untrapped electrons played a role because they were capable of following the spiral course of a perturbed field line across the containment zone. No effects due to transport derived from magnetic-perturbation gradient drifts were observed.

ACKNOWLEDGEMENT

The authors thank Professor D. W. Kerst for his suggestions and guidance. We also thank A. J. Cavallo, who developed the striped particle-collectors, for the measurements of the particle fluxes.

This work was supported by the Atomic Energy Commission.

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TABLE I

Extrema of ω_d' as a function of K and ψ
given in terms of the extrema of $B^{-1}\partial B/\partial\psi$

	ψ Region	Maximum ω_d'	Minimum ω_d'
electron	$\psi > \psi_L$	$-\frac{2K}{e} \frac{1}{B} \frac{\partial B}{\partial \psi}_{\min} + \frac{1}{L} \frac{\partial L}{\partial \psi}$	$-\frac{2K}{e} \frac{1}{B} \frac{\partial B}{\partial \psi}_{\max} - \frac{1}{L} \frac{\partial L}{\partial \psi}$
electron	$\psi < \psi_L$	$-\frac{2K}{e} \frac{1}{B} \frac{\partial B}{\partial \psi}_{\min} - \frac{1}{L} \frac{\partial L}{\partial \psi}$	$-\frac{2K}{e} \frac{1}{B} \frac{\partial B}{\partial \psi}_{\max} + \frac{1}{L} \frac{\partial L}{\partial \psi}$
ion	$\psi > \psi_L$	$\frac{2K}{e} \frac{1}{B} \frac{\partial B}{\partial \psi}_{\max} + \frac{1}{L} \frac{\partial L}{\partial \psi}$	$\frac{2K}{e} \frac{1}{B} \frac{\partial B}{\partial \psi}_{\min} - \frac{1}{L} \frac{\partial L}{\partial \psi}$
ion	$\psi < \psi_L$	$\frac{2K}{e} \frac{1}{B} \frac{\partial B}{\partial \psi}_{\max} - \frac{1}{L} \frac{\partial L}{\partial \psi}$	$\frac{2K}{e} \frac{1}{B} \frac{\partial B}{\partial \psi}_{\min} + \frac{1}{L} \frac{\partial L}{\partial \psi}$

FIGURE CAPTIONS

Figure 1: Flux plot 25 msec after the start of the main field pulse with conducting internal rings and external boundaries. Solid lines are poloidal magnetic field lines. Permanent-magnet, dipole perturbations are shown at Q and P. Note that the poloidal field lines of the pulsed containment field have soaked a distance of the order of the 2-cm skin depth into the conducting boundary but the field lines of the perturbation penetrate the boundary.

Figure 2: Curves indicating the extrema of ω_d and $\pm\omega_H$ are shown versus K for $\psi = \psi_L$. The cross-hatched area represents allowed ω_d between ω_H and $-\omega_H$ for electrons. The azimuthal angular velocity due to the $\bar{E} \times \bar{B}$ drift is given by

$$\partial\phi/\partial\psi = - (kT_e/e)/(\psi_c - \psi_s).$$

Figure 3: Azimuthal scan of floating potential showing the effects of magnetic dipole perturbations (at $\theta = 52^\circ$) applied (a) singly, or (b) as a compensating pair. All scans were made near position Q (shown on Fig. 1) on a flux surface midway between ψ_L and ψ_c . The internal ring supports were present and their effect on the floating potential is evident at $\theta = 42^\circ$.

Figure 4: Floating potential contours 3 msec after completion of the microwave heating. Supports and the magnetic dipole perturbation (nonzero $\Delta\psi$) were located at $\theta = 116^\circ$. Psi space between the ψ_s and ψ_c is shown for three cases; (a) supports present (0.25 V per contour line), (b) supports and perturbation (0.35 V per contour line), and (c) perturbation alone present (0.45 V per contour line).

Figure 5: Floating potential contours 12.5 msec after completion of the microwave heating. Supports and the magnetic dipole perturbation (nonzero $\Delta\psi$) were located at $\theta = 116^\circ$. Psi space between ψ_s and ψ_c is shown for (a) supports present (0.1 V per contour line) and (b) supports and perturbation present (0.1 V per contour line).

Azimuthal Angular Velocity — msec⁻¹