strong longitudinal field E_x was observed. The condition $\epsilon \gg (m_i/m_e)^{1/2}$ corresponds to the characteristic distance for transverse charge separation (ρ_i/ϵ) being less than the characteristic distance for longitudinal[13] charge separation $(\rho_i \rho_e)^{1/2}$. The charge separation begins at the sharp magnetic-field boundary. If the plasma is formed initially within a field or moves through an adiabatic field gradient the oscillations are not excited and the weaker condition $\epsilon \gg 1$ is still applicable.

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Explanation of the Expulsion of Impurities from Tokamak Plasmas by Neutral-Beam Injection

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Neutral-hydrogen-beam injection parallel to the plasma current (co-injection) has been observed to produce smaller concentrations of both naturally occurring and injected impurities than counter-injection. We explain this result by considering the effect of plasma rotation on radial impurity transport. This rotation effect could be the basis of an impurity-control technique for tokamak fusion reactors.

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Recent experiments on the impurity study experiment (ISX-*B*) and Princeton large torus (PLT) tokamaks have seen dramatic changes in impurity content and radial profiles during beam injection.[1,2] Beam injection parallel to the plasma current (co-injection) expels impurities while counter-injection drives them in. This points to a potentially significant method of impurity control, but it leaves us with a seeming paradox: Almost all theories predict that the effect of co-injected beam momentum is to drive impurities inward.[3-7] Those theories that predict outward flow with co-injection[8,9] must invoke a drag term with little if any physical basis to explain the result.

The purpose of the present note is to point out that the experimental results can be explained by the rapid rotation of the plasma induced by the tangentially injected beam, not by the direct beam-impurity interaction. The measured rotation speeds,[10,11] which are around 10^5 m/sec, are comparable to or greater than the thermal speed of the argon, iron, and tungsten impurities observed in the experiments. Accordingly, the impurity inertia terms, usually neglected in neoclassical theory, must be kept in the parallel momentum equation. This eventually leads to a radial impurity flux depending on the radial electric field and the main-ion parallel flow, of which the former can be eliminated by the standard relation between the electrostatic potential, the main-ion flow, and the gradients of density and temperature.

In order to make our derivation as simple as possible, we will assume that the impurity ions are in the collisional or Pfirsch-Schlüter regime.[12] (For the plasma conditions during beam injection in ISX[13] and PLT[11] and for the impurity levels given for those devices,[14,15] highly stripped argon, iron, and tungsten impurities are in the collisional regime.) We take the impurity to main-ion mass ratio to satisfy $m_I/m_i \gg 1$, so that the plasma rotation speed can be comparable to the impurity thermal speed but small relative to the hydrogen ion thermal speed. Most of our equations are simpler if stated in flux surface coordinates based on the representation of the magnetic field as $\vec{B} = I\nabla\zeta + \nabla\zeta \times \nabla\psi$.[12] The group (ψ, χ, ζ) is a right-handed coordinate system. We will denote the radius from the symmetry axis at any point on a flux surface as R. Our derivation is done in the trace-impurity limit[16] $n_I Z_I^2/n_i \ll 1$, where the presence of the impurity does not affect the main-ion flow or the electric fields. (Here, n is the density and Z is the ionic charge.) This assumption is valid for argon and iron in ISX[1] and is marginally satisfied for tungsten in PLT.[2]

In all collisionality regimes, the lowest-order flow of the main ions is given by[12]

$$\vec{v}_i = -R^2\nabla\zeta\left(\frac{1}{n_i Z_i e}\frac{\partial p_i}{\partial\psi} + \frac{\partial\Phi}{\partial\psi}\right) - \frac{I}{Z_i e} g \frac{\partial T_i}{\partial\psi}\frac{\vec{B}}{\langle B^2\rangle}, \tag{1}$$

where p_i is the pressure, $Z_i e$ is the charge, Φ is the electrostatic potential, g is a flux function whose numerical value depends on the main-ion collisionality regime, and $\langle \ldots \rangle$ denotes the flux surface average. The first term in Eq. (1) is a rigid toroidal rotation about the symmetry axis, while the second is a flow parallel to the magnetic field lines. Although Eq. (1) is quite general, we will find significant transport differences between co- and counter-injection only if the density and temperature gradient terms are comparable to the potential gradient terms. Measurements show that this is true[11] for plasmas in the plateau regime ($g = 0.5$), which is the case in the experiments.[1,2]

In order to describe the radial impurity transport, we need to solve the continuity and momentum balance equations to first order in the ordering parameter ρ/l (gyroradius divided by scale length).[12] These are

$$\nabla\cdot\vec{\Gamma}_I = 0 \tag{2}$$

$$\nabla p_I + m_I\vec{\Gamma}_I\cdot\nabla\vec{v}_I + Z_I n_I e\nabla\Phi = Z_I e\vec{\Gamma}_I\times\vec{B} + \vec{R}_I, \tag{3}$$

where $\vec{\Gamma}_I = n_I\vec{v}_I$ is the particle flux and $\vec{R}_I$ is the frictional force. We keep only the parallel component of $\vec{R}_I$ which will be written as

$$R_{\parallel I} = -(m_I Z_I^2/\tau_i) n_I (v_{\parallel I} - u_{\parallel i}), \tag{4}$$

where τ_i is the self-collision time for the main ions[17,18] and n_I is the impurity density. By using the form for the parallel flow in Eq. (1) plus the form for the collision operator between species with large mass ratio,[12] it is straightforward to derive

$$u_{\parallel i} = -\frac{I}{Z_i eB}\left[\frac{T_i}{n_i}\frac{\partial n_i}{\partial\psi} + Z_i e\frac{\partial\Phi}{\partial\psi} + \frac{\partial T_i}{\partial\psi}\left(1 - \xi + (\xi + g)\frac{B^2}{\langle B^2\rangle}\right)\right], \tag{5}$$

where ξ is a number that varies from unity when protons are collisional to 1.5 when they are in plateau regime. The terms dependent on ξ come from the thermal force portion of the frictional force.

To lowest order in ρ/l, Φ is constant on a flux surface. The next lowest-order portion, which varies on a surface, can be found in works by Hazeltine and Hinton[19] and Hinton and Rosenbluth.[20] We have found that its inclusion leads to a correction of $\lesssim 20\%$ to the flux in Eq. (9), and the additional contribution does not change sign when the plasma rotation is reversed, thus playing no role in Eq. (10) or our comparison with experiment. It is thus neglected here for brevity.

In order to have $v_{\parallel i} = O(v_{TI})$ where $v_{TI} = (2T_I/m_I)^{1/2}$, we must have $(m_I/m_i)^{1/2}\rho_{\chi i}/l$ of order one, where $\rho_{\chi i}$ is the poloidal gyroradius. This demands that we must have $Z_I \gg 1$ if we are to retain the usual relation $\rho_{\chi I}/l \ll 1$. Under this assumption, the lowest-order impurity-ion particle flux consistent with Eqs. (2) and (3) is

$$\vec{\Gamma}_I^{(0)} = -R^2\nabla\zeta(\partial\Phi/\partial\psi) n_I^{(0)} + K(\psi)\vec{B}, \tag{6}$$

where $K(\psi)$ is a flux function that must be determined and Φ is the portion constant on a flux surface.

Taking the vector product of $\vec{B}$ with the zero-order portion of Eq. (3) and employing Eqs. (5) and (6) gives the first-order, cross-field particle flux

$$\Gamma_I{}^\psi=\langle\nabla\psi\cdot\vec{\Gamma}_I{}^{(1)}\rangle=\frac{I}{Z_Ie}\left\langle\left(P_I{}^{(0)}+\frac{1}{2}m_I\frac{K^2B^2}{n_i{}^{(0)}}\right)\vec{B}\cdot\nabla\frac{1}{B^2}+\frac{1}{2}m_In_I{}^{(0)}\frac{V_E{}^2}{R^2}\vec{B}\cdot\nabla R^2\right\rangle, \tag{7}$$

where $V_E=-R\partial\Phi/\partial\psi$. This form for $\Gamma_I\psi$ shows that it is the guiding center drift, coupled with a density variation in the surface, that leads to radial transport, as we have discussed in more detail elsewhere.[17,21]

In order to solve for the density variation in the surface, we take the scalar product of $\vec{B}$ with Eq. (3) and use Eqs. (5) and (6) in order to obtain

$$\vec{B}\cdot\nabla p_I{}^{(0)}+\frac{1}{2}m_In_I{}^{(0)}\left[K^2\vec{B}\cdot\nabla\left(\frac{B}{n_I{}^{(0)}}\right)^2-\frac{V_E{}^2}{R^2}\vec{B}\cdot\nabla R^2\right]$$
$$=-\frac{m_IZ_I{}^2}{\tau_i}\left\{K(\psi)B^2+\frac{I}{Z_ie}n_I{}^{(0)}\left[\frac{T_i}{n_i}\frac{\partial n_i}{\partial\psi}+\frac{\partial T_i}{\partial\psi}\left(1-\xi+(\xi+g)\frac{B^2}{\langle B^2\rangle}\right)\right]\right\}. \tag{8}$$

Since the rotation speed is small compared to the main-ion thermal speed, the main ions have density and temperature constant on a flux surface as usual. Hence τ_i is constant. Because of the strong impurity-ion temperature equilibration, impurity temperature is also constant on a flux surface, $\vec{B}\cdot\nabla T_I{}^{(0)}=0$. Accordingly, dividing Eq. (8) by $n_I{}^{(0)}$, flux surface averaging, and using the identity $\langle\vec{B}\cdot\nabla A\rangle=0$, Eq. (8) demands

$$K(\psi)=-(I/Z_ie)[(T_i/n_i)\partial n_i/\partial\psi+(1+g)\partial T_i/\partial\psi]\langle B^2/n_I{}^{(0)}\rangle^{-1}.$$

In principle, it is possible to solve Eq. (8) in general geometry and then evaluate Eq. (7). In practice, it is much easier at this point to specialize to large aspect ratio, circular cross-section geometry, and expand in inverse aspect ratio $\epsilon\ll1$. If we take $n_I{}^{(0)}=\bar{n}_I+\tilde{n}$, where $|\tilde{n}|=O(\epsilon\bar{n}_I)$ and $\langle\tilde{n}\rangle=0$, we can solve Eq. (8) for $\tilde{n}=A\sin\chi+C\cos\chi$, where

$$A=-(m_iZ_I{}^2/\tau_iZ_ieB_{\zeta0})\bar{n}_Iq^2R_0[\delta^2+(1-y)^2]^{-1}$$
$$\times\{(2-y+z)(T_i/n_i)\partial n_i/\partial r+[2(1-\xi)+y(2\xi+g-1)+z(1+g)\partial T_i/\partial r]\},$$

$$C=\epsilon\bar{n}_I[\delta^2+(1-y)^2]^{-1}\{(y+z)(1-y)-2\delta^2[(T_i/n_i)\partial n_i/\partial r+(1-\xi)\partial T_i/\partial r][(T_i/n_i)\partial n_i/\partial r+(1+g)\partial T_i/\partial r]^{-1}\},$$

$$\delta=(m_iZ_I{}^2/\tau_i)qR_0V_\parallel/T_I,\quad V_\parallel=-(Z_ieB_{\chi0})^{-1}[(T_i/n_i)\partial n_i/\partial r+(1+g)\partial T_i/\partial r],\quad y=m_IV_\parallel{}^2/T_I,\quad z=m_IV_{E0}{}^2/T_I.$$

With use of the form for $\tilde{n}$ in Eq. (7), the cross-field flux $\Gamma_{Ir}=\Gamma_I{}^\psi/B_{\chi0}R_0$ is

$$\Gamma_{Ir}=\frac{m_i}{\tau_i}\frac{2q^2}{e^2B_0{}^2}\frac{Z_I}{Z_i}\bar{n}_I\frac{1-(1/2)(y-z)}{\delta^2+(1-y)^2}\{[1-\tfrac{1}{2}(y-z)](T_i/n_i)\partial n_i/\partial r$$
$$+[1-\xi+\tfrac{1}{2}y(2\xi+g-1)+\tfrac{1}{2}z(1+g)]\partial T_i/\partial r\}. \tag{9}$$

The solution in Eq. (9) is valid only as long as y is not too near unity, so that $\tilde{n}=O(\epsilon\bar{n}_I)$ remains true.

Equation (9) shows that the parallel impurity flow caused by frictional coupling (i.e., $V_\parallel$ and y) and the toroidal impurity flow due to the radial electric field (V_{E0} and z) provide quite different effects. Parallel impurity flow can cause outward impurity convection if it is large enough; toroidal flow simply enhances inward convection.

If we restate Eq. (9) in terms of the experimentally measured rotation speed $U=V_\parallel+V_{E0}$, we can easily show that there is a difference in radial transport for co-rotation ($U>0$) and counter-rotation ($U<0$). It is observed[11] that $U\to-U$ when the beam direction is changed if the magnitude of the beam momentum input is the same. Under the assumption that $n_i(r)$ and $T_i(r)$ remain the same, the change in flux is

$$(\Gamma_{Ir})_{\rm co}-(\Gamma_{Ir})_{\rm ctr}=-\frac{m_i}{\tau_i}\frac{2q^2}{e^2B_0{}^2}\frac{Z_I}{Z_i}\frac{2m_I|U|V_\parallel}{T_I}\frac{\bar{n}_I}{\delta^2+(1-y)^2}$$
$$\times\left\{\left(2+\frac{m_IU^2}{T_I}\right)\frac{T_i}{n_i}\frac{\partial n_i}{\partial r}+\left[2-\xi+g+\frac{m_I}{T_I}\left(U^2(1+g)+V_\parallel{}^2(\xi+g)\right)\right]\frac{\partial T_i}{\partial r}\right\}. \tag{10}$$

Equations (9) and (10) make predictions that agree with experiment. First, for density and temperature gradients of the usual sign, Eq. (10) predicts a greater influx of impurities for counter-injection, as is seen in experiments.[1,2] Second, if we use representative values for PLT,[2] the difference in the radial convection velocities $v_{Ir} = \Gamma_{Ir}/\bar{n}_I$ is in the range 0–10 m/sec, compared to the actual values of 0–4 m/sec. Third, if we consider that diffusion as well as convection will be taking place in the plasma, then the ratios of coefficients in Table I show that transport in discharges with only co-injection and those with both co- and counter-injection should be much more similar to each other than either is to counter-injected discharges. This is indeed seen in PLT.[2]

The theory makes specific predictions that can be tested experimentally. The sign of the impurity flux should be independent of the sign of the toroidal field B_ζ. In addition, for fixed q, the magnitude of the outward flux should scale as $|B_\zeta|^{-2}$. (Consequently, ISX should see a much larger effect than PLT.) Finally, the result in Eq. (9) is quite small near $y = z + 2$. A test with impurities of different mass might be used to detect this feature of the result. The predicted impurity density variation $\tilde{n}$ can also be used to check the theory experimentally. Except very near to $y = 1$, this variation is primarily $\cos\chi$, which is unlike the $\sin\chi$ variations seen previously in the absence of beam injection.[22]

In conclusion, we can explain the observed difference of behavior of argon, iron, and tungsten under co- and counter-injection by considering the effect of plasma rotation on impurity transport. Using this effect to remove impurities from the center of a reactor offers the possibility for creating an ideal reactor plasma, which is clean in the center but dirty and radiation dominated at the edge. Since no direct momentum to the impurity ions is required, the rotation can be created by any means (e.g., neutral beams or electromagnetic waves) which can impart momentum to the main ions.

TABLE I. Ratios of coefficients for PLT (Ref. 1).

	Coefficient of $n_i^{-1}\,\partial n_i/\partial r$	Coefficient of $T_i^{-1}\,\partial T_i/\partial r$
Conventional theory ($y = z = 0$)	1	0
Co-injection	0.16	−0.56
Counter-injection	4	10
Both[a]	0.013	0.22

[a]Assumes T_i increases by a factor of 1.5 from co- or counter-injection cases, since the power input from one beam is approximately equal to the Ohmic input.

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Energy Partition in CO_2-Laser-Irradiated Microballoons

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A self-consistent study of the partition of energy absorbed in nanosecond single-CO_2-laser-beam irradiation of glass microballoons is reported. Through interferometric inference of shell heating, and quantitative fast-ion spectrometry, it is shown that a major fraction of the absorbed energy is converted to fast-ion expansion and only 25% contributes to thermal heating of the target.

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Studies of nanosecond CO_2-laser irradiation of plasmas at intensities of $>10^{13}$ W cm^{-2} have shown that absorption primarily occurs through collisionless processes, particularly resonance absorption,[1] which generate superthermal electrons[2] that are confined by the ambipolar field they establish with the ions to form a hot low-density corona. These collisionless electrons will in general follow nonradial paths[3] in this coronal plasma and drive its expansion. On each encounter with the plasma sheath, an electron loses a small fraction $\sim(Zm_e/m_i)^{1/2}$ of its energy to fast-ion expansion[4] and thus many reflections are required to convert a sizable fraction of the absorbed energy to fast-ion expansion.

In nanosecond irradiation of spherical targets, the sheath will expand with a characteristic speed $c_h \sim (ZkT_h/m_i)^{1/2}$, where T_h is the hot-electron temperature. Consequently, for long laser pulses, it can expand to dimensions such that the geometric cross section of the target is small.[5] Only those electrons whose orbits do intersect the dense cool plasma may contribute to heating of the target. Thus, the fraction of energy converted to fast-ion expansion depends upon a number of factors including the value of T_h, the laser pulse duration, and target geometry. Various estimates of this loss ranging from 9% to 90% have been inferred through indirect means.[6]

The experiments reported here are the first self-consistent measurements of the energy partition in a CO_2-laser-produced plasma. Independent estimates were obtained for the total energy absorbed, the energy lost to fast ions, and the energy deposited into the dense target. The targets used were of three types: empty glass microballoons of 150 and 220 μm diameter with a wall thickness 1.5 μm; empty glass microballoons of 200 μm diameter, coated with 20-μm-thick $(CH_2)_n$; and solid glass spheres of 220 μm diameter. These targets were irradiated by single 20-J, 1.4-ns [full width at half maximum (FWHM)] pulses from the COCO-II laser system. The $f/2.5$ center-focused beam had a half-energy diameter of 110 μm.

An estimate of the total thermal energy deposited in the target was inferred from interferograms taken at various times during and following irradiation of the target with the aid of a synchronized 0.53-μm, 70-ps probe pulse and a folded wave front interferometer. The plasma in the focal region shows strong density-profile steepening, indicative of ponderomotive and superthermal pressure effects at the critical density.[7] However, away from the focal region the electron density resembles that expected for a simple isothermal expansion, Fig. 1(a), from which a characteristic scale length, L, may be obtained.