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HI ATOMIC BEAM SOURCE: A FEASIBILITY STUDY

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ABSTRACT

The possible application of the stable atomic hydrogen (H_I) in intense polarized atom or ion beam sources is discussed. The proposed scheme is based on flipping the electron spin of stored atoms by microwaves; this results in the magnetic acceleration and focusing of the flipped atoms. The main advantages of such beams are: accumulation and pulsing with time structure convenient for injection into synchrotrons via Linac injectors; high density due to the slow velocity and high initial density; and monochromaticity which allows atomic beam transport with focusing elements having no large chromatic errors. The possible problems and limitations of such an extraction scheme are: depolarization and slow spin diffusion due to spin-exchange collisions in the dense gas phase, and loss of density due to recombination of the flipped atoms on the walls before reaching the extraction aperture and accelerating field gradient.

The results of the first tests of our apparatus are presented. The apparatus, based on a powerful dilution refrigerator, is described briefly.

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INTRODUCTION

Soon after the experimental discovery of the stable gaseous state of atomic hydrogen [1], it was proposed that the atoms stored in the magnetic "bottle" could be extracted by flipping their electron spin by resonant microwaves [2]. It was suggested that the magnetic acceleration and focusing by the solenoid field would result in a polarized, high-density, monochromatic atomic beam, which could be useful in polarized ion sources, polarized atomic beam ("jet") targets, and atomic beam storage rings [2].

The use of stable atomic hydrogen in ion sources would provide the inherent capability of storing the atoms in between the output pulses. The resonant microwave pulse extraction could take full benefit of this capability, which has the potential of injecting polarized protons into synchrotrons with intensity equal to that of unpolarized beams.

In d.c. beam sources, however, the stable atomic hydrogen might offer little improvement compared with the best (moderately cold) classical polarized atom beam sources, because the losses in the cooling and stabilization might well offset the gains due to focusing and somewhat lower velocity.

Although our aim here is not to review extensively the classical ion sources, a few comments will be made on them in order to clarify the above statements and to understand the scope of this work.

The cooling of the discharge or nozzle of an atomic beam source has given an improvement proportional to $T^{-1/2}$ instead of the expected $T^{-3/2}$ [3]. This is partly due to the slower flow through the nozzle at lower temperature, which has not been properly considered in the calculation of the improvement factor. The main problem of these sources is their chromaticity. It has been speculated that the slower velocity of the cooled atoms would permit a larger angle of acceptance in the long hexapole magnet channel which is customarily used in these sources for state selection. The larger angle of acceptance, however, results in a larger chromatic aberration and virtually no gain in the density at the focal point.

The use of short hexapole lenses is limited by the available maximum radial field gradient. The use of short superconducting solenoid lenses could possibly reduce the chromaticity problems in cooled classical atom beam sources [2,4]. A cold atom source could be placed very near to a short high-field solenoid lens, the main purpose of which would be to focus the atoms with electron spin up, and

defocus (possibly store) the stable (spin down) species. This scheme was first suggested in a discussion at the 1980 Lausanne Symposium. The scheme differs from a hexapole polarized atom source only by the fact that a superconducting solenoid would allow $\mu B \approx E_{kin}$ to be satisfied, where μ is the electron magnetic moment and E_{kin} the maximum useful kinetic energy of the atoms emerging from the discharge nozzle or accommodator. A strong divergent field of a short solenoid can then have a good focusing effect on the c- and d-state atoms, whereas atoms in the a- and b-states are attracted towards the highest field near the winding of the solenoid, and become either recombined on solid H_2 coated walls, or cooled and stored in a cell the walls of which are coated with superfluid 4He . These schemes are

Fig. 1: Cold atom beam source with solenoid spin state separating magnet. The accommodator temperature should be 3 K or higher.

shown in Fig. 1; the cryopump panels play an important role in any cold source because no mechanical or diffusion pump can handle the gas load of the projected high intensity sources.

At the 1981 Workshop on Polarized Proton Ion Sources the discussion on cold atomic hydrogen concentrated on direct ionization in the storage cell by a neutral Cs beam [5], or on its use as a polarized target or atom beam source with an unspecified extraction mechanism [6]. A scheme was also discussed where an unpolarized atom jet in a high-field area would be cooled to 0.5 K in a helium-coated accommodator, before an RF transition for the nuclear polarization. It is, however, easy to show that this proposed scheme would not work because the cooling to low temperature in high field will inevitably lead to electron spin polarization; the escaping atomic beam would be highly divergent due to the divergence of the solenoid field at the end. Also,

the d.c. mode of operation in this scheme would give no or little advantage in synchrotron ion sources, because most of the atom flux would be lost due to the duty cycle of the injector.

In the d.c. beam sources the basic limitation is the discharge atom source; because losses in the subsequent cooling are the greater the lower the temperature, little advantage may be expected of the "ultra-cold" sources compared to the moderately cold ones. In particular, little benefit can be expected from the larger angle of acceptance in the first focusing element, because of the chromaticity of the source.

In this paper the direct ionization in the stabilization cell [5] will not be discussed; we shall concentrate on the microwave flip extraction of double polarized atoms in view of multiple applications also without ionization. Our project has been briefly described in recent publications [7].

EPR SATURATION OF HI

The Electron Paramagnetic Resonance (EPR) saturation of gaseous HI cannot be calculated in analogy with solids, because even at low temperature the atoms move fast enough for perfect motional narrowing. In analogy with resonance crossing in polarized beams, we calculate here the resonance strength required for adiabatic reversal of the electron spin when an atom crosses the surface where $|\vec{H}| = H_0 = (g\beta)^{-1}h\nu_0$ (ν_0 is the microwave frequency, $|\vec{H}|$ the value of the magnetic field, and g the g -factor of the particular hyperfine transition).

For the crossing to be adiabatic it is required that during the passage through zero effective field, the spin makes many Larmor precession cycles. The apparent resonance width only depends on the microwave field H_1 perpendicular to the main field:

$$\Delta\nu = g\beta H_1/h, \quad (1)$$

and the crossing lasts

$$\Delta t = 2H_1(\vec{v} \cdot \nabla |\vec{H}|)^{-1} \quad (2)$$

for an atom with velocity $\vec{v}$. The spin will make many precession cycles if $\Delta t \gg (\Delta\nu)^{-1}$; hence

$$H_1^2 \gg h(\vec{v} \cdot \nabla |\vec{H}|)(g\beta)^{-1}, \quad (3)$$

where $\langle \rangle$ denotes thermal average over the Maxwell-Boltzmann distribution. At 0.3 K temperature and 5×10^4 Oe field the requirement is $H_1 \gg 0.1$ Oe, which is not very difficult to achieve even in large cavities.

In the case of a large multimode cavity, the requirement of microwave power P_{rf} becomes

$$P_{rf} \approx \mu_0 H_1^2 V v / (2Q) \gg \mu_0 H_0^2 V \langle v \rangle \Delta H / (2L H_0 Q), \quad (4)$$

where V and Q are the effective volume and quality factor of the cavity, $\Delta H / (H_0 L)$ the average logarithmic field gradient in the cavity, and $\langle v \rangle$ the mean thermal velocity. In our case $V = 2 \text{ cm}^3$, $Q \approx 10$, $\Delta H / L H_0 = 10^{-4} \text{ cm}^{-1}$, $L \approx 1 \text{ cm}$, $H_0 = 5 \times 10^4$ Oe, giving $P_{rf} \gg 1 \text{ W}$ for complete spin flip of all atoms crossing the resonance surface. The time required for turning all spins would then be about $L \langle v \rangle^{-1} \approx 0.2 \text{ ms}$.

The fate of the sample under continuous microwave irradiation is rather complex. An atom with flipped spin may

- flip back by a second passage through resonance;
- be accelerated by the field gradient and ejected through the filling tube;
- hit another atom in the gas phase with opposite electron spin; these collisions may or may not be associated with electron spin exchange and may or may not depolarize the nuclear spin;
- stick to the wall and recombine;
- stick to the wall and escape with or without wall depolarization.

The collisional processes, which limit the drift and diffusion of the atoms and cause depolarization, will be discussed in the next section.

At saturating microwave power the initial velocity distribution of the flipped atoms is likely to be rather isotropic. At reduced power the atoms crossing the resonant surface at a skew angle will have higher probability of a spin flip, which may result in the initial velocity peaking at right angles with respect to the axis of the beam line. This may be improved by applying a suitable field gradient along the axis and sweeping the microwave frequency in such a way that the resonant surface will follow, at the mean thermal velocity, the atoms which have their velocity vector pointing in the direction of the extraction aperture. A frequency sweep somewhat faster than the mean thermal velocity could possibly increase the effective volume in which the spin flip occurs, and therefore further increase the number of extracted atoms and their initial directionality. Solid-state pulsed

microwave sources are available with 'chirp' naturally occurring during the pulse.

SPIN DIFFUSION AND NUCLEAR DEPOLARIZATION

In the following we shall discuss the phenomenology of the gaseous atomic hydrogen consisting of atoms in the two lowest (a- and b-) states and in one of the two higher (c- or d-) states. The partial densities n_a and n_b are not necessarily assumed to be equal; the higher state concentration n_i ($i = c$ or d) is supposed to be small in comparison with the total density n .

In a relatively dense gas phase the atoms with flipped electron spin will diffuse away from the resonance surface at a speed determined by the classical equation

$$\partial n_i / \partial t = D \nabla^2 n_i, \quad (5)$$

where the diffusion constant D is

$$D = t_c kT/m; \quad (6)$$

here we approximate t_c , the inverse of the mean collision frequency, by the mean free time τ : $t_c \approx \tau = \langle \lambda/v \rangle = \lambda v / (2m/\pi kT)$. The mean free path $\lambda = (n\sigma_c)^{-1}$ is determined by the singlet collision cross-section of the flipped species against the lower state background gas; we shall return to this point later in this section. The resulting approximate diffusion coefficient is

$$D = \langle v \rangle / (2n\sigma_c). \quad (7)$$

The flipped atoms which reach an area of large field gradient will drift with a velocity

$$v_z = \Omega \nabla |H|, \quad (8)$$

where the mobility Ω is given by

$$\Omega = \mu_B t_c / m \approx \mu_B \lambda / (m \langle v \rangle). \quad (9)$$

We may note that if the collision cross-section σ_c does not depend significantly on the relative particle velocities, the diffusion constant

and mobility have an accurate relation $D/\Omega = kT/\mu_B$.

The collisions which limit the diffusive and drift motions may also depolarize the nuclear spins of the flipped atoms. If the collision cross-section is not much larger than the spin-exchange cross-section (i.e. the spin flip probability is near the maximum of 1/2 in collisions), the equations (5) and (8) above should be corrected for the state non-conserving exchange collisions.

Many of the completely spin-polarized gas properties have been calculated already. In particular the electron spin depolarization due to exchange and dipolar interactions has been calculated for low temperature and high field by Kagan et al. [8]. The electron spin exchange there arises from the small admixture of electron spin up in the a-state atom. This inelastic cross-section is small but not negligible from the point of view of recombination at high density. In our case, where the singlet pairs are supposed to exist in the gas phase, these exchange terms, however, can be neglected. On the other hand, the spin-orbit interaction, which is small for the repulsive triplet potential, may become notable for the singlet potential, leading to spin-orbit depolarization, which is the quantum analogue of Thomas precession. We note that this depolarization also happens in collisions with helium atoms which may be present.

The transport coefficients in the completely spin-polarized gas have been calculated by Lhuillier and Laloë [9]. The nuclear spin diffusion is shown to be connected with heat conduction. Unfortunately these results are not applicable to our case which involves the added complexity of the higher hyperfine states in the gas.

Earlier theoretical work involving singlet pair collisions has been restricted to the completely depolarized gas at zero field and room temperature [10,11] or at zero field and moderate [12] or low [13] temperature. The electron spin exchange in singlet pair collisions of hydrogen atoms in high field can be determined from the singlet potential which crosses zero at an interatomic distance of about 0.4 Å. This de Broglie wavelength corresponds to a mean thermal velocity around room temperature. Although the attractive part of the potential extends about ten times further, we might assume for simplicity that the s-wave scattering dominates the $\ell \geq 1$ terms for singlet pair scattering. Kagan, Vartanyants and Shlyapnikov [8] have calculated the s-wave scattering length $a_s = 0.33$ Å using the singlet potential of Kolos and Wolniewicz [14]. Kagan et al. point out that, for triplet pairs below 1 K, the scattering can be described to a good approximation by the $\ell = 0$ and 1 partial waves, in good agreement

with the phase shifts of Ref. [9]. For simplicity we assume that $\ell \geq 1$ partial waves for singlet pairs can be neglected, which leads to a velocity-independent cross-section $\sigma_s = \sigma_s^0 = 4\pi a_s^2$ and, furthermore, because the attractive potential brings the atom pair rather close together, we assume equal probability for the electronic spin-exchange and spin-conservation. We might thus use the exchange cross-section

$$\sigma_s^{\text{ex}} = 2\pi a_s^2 = 0.683 \times 10^{-16} \text{ cm}^2. \quad (10)$$

We may compare this with the zero-field value $\sigma^{\text{ex}} = (0.61 \pm 0.03) \text{ \AA}^2$ determined from the longitudinal relaxation rate at 1 K [15]. The agreement must be considered accidental because at zero field the calculation of σ^{ex} involves phase shifts arising from both singlet and triplet potentials [12,13]. At high field and low temperature, ignoring the wall collisions (i.e. assuming a high density and that the mean free path is small in comparison with the size of the cell), the spin-exchange collision rate becomes

$$\tau_{\text{ex}}^{-1} = \langle \sigma_s^{\text{ex}} v \rangle n \cong \sigma_s^{\text{ex}} \langle v \rangle n. \quad (11)$$

Here we have made the assumption that at any time the concentration of the flipped atoms is small in comparison with that of the a- and b-state atoms, which is plausible because both the recombination and the magnetic field expulsion will provide efficient sinks for the upper hyperfine state atoms.

We note that in {a,d} and {b,c} collisions the electron spin exchange will result in the same state atoms thus conserving the total hyperfine energy and causing no nuclear spin depolarization, but in collisions {a,c} and {b,d} we get

$$\{a,c\} \leftrightarrow \{b,d\} - A/2, \quad (12)$$

where A is the hyperfine energy. Although the inelasticity in the collisions is very small, it may also slightly modify the cross-sections for nuclear depolarization in the d- or c-state atoms. Neglecting these (possibly) small corrections, the d- and c-state atoms will convert into each other with rates

$$\tau^{-1}(d) = \sigma^{\text{ex}} \langle v \rangle n_b \quad (13)$$

$$\tau^{-1}(c) = \sigma^{\text{ex}} \langle v \rangle n_a. \quad (14)$$

We note that because the b-state is more stable and becomes naturally enriched over the a-state, it might be much more favourable to extract the c-state atoms with microwave transitions b→c. From the point of view of beam handling, however, it would be convenient to deal with atoms in the pure d-state, which does not depolarize at low field. The transition b→d ($\Delta l = 2$) is possible by a 2-photon process [16] but requires a high microwave power.

EXPERIMENTAL APPARATUS AND RESULTS

Figure 2 shows schematically the CERN apparatus for the

Fig. 2: The CERN hydrogen stabilization apparatus. The thermal shields and vacuum jackets are not shown for clarity.

experimental studies of the stable atomic hydrogen. The horizontal dilution refrigerator [17] has a maximum cooling power of 350 mW at 0.5 K, 80 mW at 0.3 K, and 30 mW at 0.2 K. In the present studies about 50 mW constant heating power was applied to the helium mixture in the mixing chamber in order to keep the cell temperature at about 300 mK; the resulting temperature stability was somewhat insufficient for the EPR measurements owing to the oscillatory behaviour of the refrigerator in this temperature range.

The hydrogen loading and stabilization assembly, which is an independent insert, is mounted in the direct access tube of the dilution refrigerator. This assembly carries the PTFE loading tube of 8 mm diameter and 70 cm length from the back plate to the mixing chamber. The PTFE tube is shielded by a corrugated stainless-steel tube; the 4He film creeps in between these tubes and causes their

thermalization. The transition from 30 K to 0.5 K is made as short as possible by applying power to electric heaters on the loading tube.

The cell area of the dilution refrigerator is surrounded by a warm-bore superconducting solenoid, which can be trimmed to give a 5 T field of 10^{-4} homogeneity in a volume of about 1 cm^3 , or a 6 T peak field, by adjusting the currents in its three independently controllable coils. The solenoid dewar can be positioned after suitable field shaping; a ^3He NMR probe in the mixing chamber near the cell facilitates these operations.

Figure 3 shows the sample cell area in greater detail. The

Fig. 3: The accommodator and cell in the mixing chamber.

stabilization cell can be exposed to microwaves through a thin cylindrical mylar window. A carbon resistor thermometer in the microwave cavity monitors the microwave field strength; this bolometer was used for the observation of the EPR transition signals.

The first tests were done with a somewhat simpler cell with a recombination bolometer directly mounted in the cell. Figure 4 a shows a typical signal of the cell bolometer current, which is measured by applying a constant voltage bias over the series circuit of the bolometer and a reference resistor, and recording the voltage of the reference resistor.

The bolometer response was calibrated by comparison with a calorimetric signal (Fig. 4 b) from another carbon resistor thermometer attached to the cell body; the absolute calibration was obtained by feeding electric pulses of known energy to a heater on the cell while recording the thermometer reading.

The hydrogen discharge tube, made of quartz, is mounted in a coaxial cavity fed by a magnetron at 2.45 GHz; about 50 W power was

Fig. 4: Typical recombination signals of the bolometer and calorimeter for a high-density sample.

used to dissociate the H_2 at $2 \mu\text{mol/s}$ flow rate, which corresponds to normal operating conditions in our case. The discharge pressure is then about 0.6 Torr. The normal loading speed, as determined by recombination bolometry, is 10^{16} s^{-1} at 5 T field and 0.3 K helium temperature; a higher discharge pressure resulted in a higher flux and impractically short cell-filling times. The normal flux causes about 1 mW heating in the mixing chamber. The transmission in the PTFE loading tube degraded in the time scale of a few hours under continuous loading. This is believed to be due to formation of snow-like solid hydrogen in the loading tube filled by low pressure ^4He owing to the climbing and evaporation of the superfluid film into the teflon tube. The loading speed could be recovered by a warm-up cycle above 30 K.

Our best yield of atoms in the cell is more than two orders of magnitude below the molar flow of hydrogen into the discharge; we tend to explain this by a poor dissociation ratio (estimated at about 10%) and recombination losses in those parts of the PTFE tube which become coated by molecules degassing out of the hydrogen pipework after an LN_2 -temperature charcoal purifier.

The maximum sample density accurately determined from the calorimeter signal is $5 \times 10^{16} \text{ cm}^{-3}$. By extrapolating outside the calibration range, bolometric signals indicate maximum densities around 10^{17} cm^{-3} , which approaches the theoretical maximum at 5 T.

The EPR measurements were performed by slowly sweeping the magnetic field over the resonance frequency, recording the differential thermal signal from the cavity and mixing chamber resistors. The microwave absorption by the sample caused heating of the cavity relative to the mixing chamber, with polarity opposite to what was

expected if a bolometric signal could have been seen. In our thermally linked geometry, where only the Kapitza resistance provides thermal isolation between the various areas, the thermal and bolometric signals partly compensate each other, making quantitative estimates difficult.

Figures 5 a and b show the direct and differential thermal signals

Fig. 5: EPR-induced recombination signals for a-d (a) and b-c (b) transitions. Note the more sensitive scale in (a).

for the a→d and b→d transitions. The direct signals demonstrate the drift and oscillation of the mixing chamber temperature, which nearly prevented the observations of the weak thermal signal of the EPR induced recombination. The microwave power level for these records was about 40 nW, most of which was absorbed in the cavity bolometer, and the field was swept at a speed of 15 G/min. The stability of the 140 GHz carcinotron source was about 5 MHz, determined by down-conversion in a harmonic mixer.

Figure 6 shows the time evolution of the signals corresponding to the two transitions at 0.3 K temperature. The a→d signal goes down more slowly than one would expect from the theory. We suggest that this is due to a-b relaxation. The density of the samples, after each series of EPR scans, was determined by triggering the recombination bolometer; the measured density values are indicated in Fig. 6 for each sample.

Although our sample cell geometry was designed for the possible direct observation of the microwave flip extraction of polarized atoms into the loading tube, the thermal fluctuations prevented sufficiently accurate quantitative measurements. The only conclusion at this stage

Fig. 6: Time evolution of the EPR signal strengths. The vertical scale shows the number of spins flipped during the scan of the magnetic field through the resonance. The sample densities, determined from recombination heat right after the last scan of each series, are: $\blacksquare$: 10^{16} cm^{-3} ; $+$: $5 \times 10^{16} \text{ cm}^{-3}$; $\blacktriangle$: $5 \times 10^{16} \text{ cm}^{-3}$.

is that resonant microwave power up to 1 mW in our cavity does not cause uncontrolled recombination, and that the spin flips at low microwave power are detected by the thermal signal due to recombination.

ATOMIC BEAM EXTRACTION AND OPTICS

The planned modification of our apparatus is shown schematically in Fig. 7. We aim at extracting and detecting an intense atomic beam by microwave-induced spin flips. In this design the heated PTFE loading tube is carried only to the level of the still, where all surfaces are coated by the superfluid film. The emerging 30 K atoms are cooled in contact with baffles around 0.5 K temperature before entering the high-field storage tube, which should be stabilized at a temperature somewhat below 0.3 K.

Fig. 7: Planned modification of the CERN atomic beam source. Note the increased cold volumes, and use of the cooling available on the still and heat exchanger walls.

A microwave reflector in the cold tube prevents direct feed of the radiation to the loading port, and reduces also the direct flight of the hot atoms through the cold tube. The microwave pulses from a feed horn are focused by a parabolic mirror to the extraction orifice, which is located at a relatively low field in order to reduce the thermal leakage out of the cell. The extraction orifice may be surrounded by a small superconducting coil with opposite field; the purpose of this is to further reduce the leakage, and to produce a magnetic 'funnel' near the focus of the solenoid end.

The open geometry of the apparatus requires well-localized ^4He film burners and efficient vapour-condenser baffles, in order to avoid excessive heat leaks to the dilution refrigerator, and to prevent the contamination of the beam vacuum with ^4He . Figure 7 shows the planned configuration of the baffles, which permit no direct passage. The baffles have a helical structure to allow easy assembly of the inner and outer pieces.

The magnetic acceleration out of the storage solenoid field B_s will result in the shrinkage of the axial velocity distribution, and in the radial focusing of the beam due to the divergence of the solenoid field. The final axial velocity becomes $v_f = \sqrt{v_m^2 + v_0^2}$, where $v_m = \sqrt{2\mu B_s/m}$ is the velocity due to the magnetic potential and v_0 is the initial axial velocity. The acceleration will shrink the axial velocity distribution by

$$\Delta v_f \cong \Delta v_0 / R,$$

where the shrinkage factor $R = \sqrt{(2\mu B_s/kT)}$. At $T = 0.3$ K and $B_s = 5$ T we get $R \approx 5$; the above estimate, however, does not properly take into account the change in the shape of the velocity distribution after the magnetic acceleration, and is therefore somewhat pessimistic.

Preliminary results from track-tracing calculations [18] suggest that the acceptance angle of the initial velocity of the flipped-spin atoms can be as large as 60° off the axis, for a 20 mm extraction orifice. This demonstrates convincingly the advantages of the monochromaticity and strong focusing due to magnetic field divergence.

A convenient way of discussing the optics of the rather monochromatic atomic beams is to define a focal length for the magnetic elements of the assembly, the main components of which are the half-solenoid of the storage cell itself, and possible short solenoid lenses.

The field of the main solenoid can be approximated by that produced by a thin solenoid. We define the focal length

$$L_f = r(v_z/v_r)_\infty, \quad (15)$$

where the radial and axial speeds v_r and v_z are taken far from the optical element (at zero field), and the thickness of the element is assumed small, i.e. the atom moves at constant r in the field. We get for the half-solenoid of the source

$$L_f^{hs} = 2a[1 + (v_0/v_m)^2]^{1/2} \mathfrak{S}_{hs}^{-1}, \quad (16)$$

where v_0 is the initial thermal velocity of the atom, and $v_m = (2\mu B/m)^{1/2}$ is the magnetic velocity due to the magnetic potential; $\mathfrak{S}_{hs}$ is a definite integral

$$\mathfrak{S}_{hs} = \int_0^\infty (a^2/r) \{(\nabla \vec{B})_r/B_0\} \{v_m/v_z(z)\} (dz/a), \quad (17)$$

which depends on v_0 . The approximations made in the above are good to about $r = 0.2a$, where a is the mean radius of the solenoid.

The numeric integration for a thin solenoid approximating our magnet gives $L_f^{hs} = 100$ mm for " T_0 " ($=mv_0^2/2k$) = 0.3 K and central field $B_0 = 5$ T, in excellent agreement with track-tracing calculations [18].

For the passage of parallel "atom rays" through a short solenoid magnet of radius a , we get a focal length

$$L_f^{ss} = 2a(B_s/B_0)\mathfrak{G}_{ss}^{-1}, \quad (18)$$

where

$$\mathfrak{G}_{ss} = 2 \int_0^\infty (a^2/r) \{(\nabla B)_r/B_0\} [1-B_z/B_s]^{-1/2} (dz/a). \quad (19)$$

In the above integral B_s is the source field giving the initial velocity $v_z(-\infty) = (2\mu B_s/m)^{1/2}$ far before the coil loop, and B_0 is the central field of the loop. For $a = 75$ mm and $B_0/B_s = 0.9$ we get $L_f^{ss} = 500$ mm. Much stronger focusing is possible by suppressing the field of the coil far from it; this can best be done by using gradient coil pairs with an axial gap of 0.5 to 1.0 coil radii. These preliminary results look very promising in view of the use of large-aperture thin-lens optics for transporting the monochromatic atomic beams.

CONCLUSIONS

An atomic hydrogen beam, spin-flip extracted from a stabilization cell, is predicted to have the following characteristics:

- 1) nuclear polarization near 100%;
- 2) monochromaticity: $\Delta v/v \approx 0.05$ (axial beam temperature about 10 mK);
- 3) mean velocity ≈ 200 m/s;
- 4) possible beam density in pulsed operation $\geq 10^{13}$ cm $^{-3}$.

The above values are based on the present CERN hydrogen stabilization apparatus; they could be improved by better cooling and a more powerful discharge atom source. The largest refrigerator at CERN might allow d.c. operation at a maximum flux of about 10^{22} s $^{-1}$, which approaches the range of interest in plasma heating by a polarized neutral beam.

The monochromaticity of the beam could make possible the storage and accumulation of the atomic beam pulses in a magnetic ring, where acceleration, deceleration, and further cooling of the beam might be possible.

The physics potential of these cold atomic beams, in addition to the polarized ion sources and polarized atom beam targets, includes also the high-accuracy measurements of the parity violation in the hydrogen atom, provided that a suitable way of exciting to the 2s state can be developed.

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