

Letters to the Editor

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Temporal Growth of Current between Parallel Plates

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IN the theory of current growth between parallel plates in a gas (the Townsend discharge), an important case is that in which there are two secondary processes in operation at the cathode, *viz.*, electron generation by photons (δ process) and by positive ions (γ process). A well-known problem concerning this system is the calculation of the cathode electron current at any time (in absence of space charge distortion) produced by exposing the cathode to a constant external illumination from time zero onwards, the gas being initially free from charged particles. If αd (where α is the first Townsend coefficient and d the plate separation) has a typical experimental value of the order of ten, the exact solution, either in the form in which I originally gave it,¹ or in the equivalent form in which I have since given it,² is, owing to its absolute accuracy, rather cumbersome if applied rigorously up to large multiples of the electron transit time. Writers have therefore directed more attention to my approximate formula,¹ and particularly to the special case in which (by taking one of the secondary coefficients, say γ , to be zero) it is applied to the simpler problem in which there is only one secondary process. Auer,³ in a study of this simpler problem, has recommended a slight modification of my approximate formula for it; but, for reasons which I have pointed out,² the modification which he recommended is not desirable.

Auer has recently proposed⁴ two new formulas (still for the special case of γ zero). One is proposed as a new equivalent form of the exact solution, and the other as an approximation to it. He obtains these formulas by a procedure which is interesting, but which contains a mathematical fallacy. For typical values of αd of order ten, they are in consequence quite wrong, except in the first electron transit time. After only about three electron transit times they are not even of the right order of magnitude. Suppose, for example, that, with αd about ten, we consider a slightly overvolted gap, (or one which is slightly undervolted

and in which the electron current at the cathode attains, in reality, an amplification of, say, a hundred, after an infinite time). In all such cases the amplification after only three electron transit times will, in reality, be only about four; but Auer's new formulas predict an amplification or order ten to the power of eight. If the over (or under) voltage is increased, the predictions remain equally unrealistic.

The incorrectness of Auer's new formulas is due to an erroneous mathematical argument, not employed in his previous paper. In setting up Eq. (4) of his new paper, he has tacitly assumed that the quantity which he calls h has to be continuous at the positive integral values of ξ . But actually it is f that has to be continuous. As ξ passes through an integral value, the n in his Eq. (2) changes suddenly by unity, and h must change suddenly to prevent f from changing. Applying this at $\xi=1$, and taking αd (and hence also Auer's q) to be of order ten, it will be found that as ξ passes through unity, h falls suddenly from a value unity to an extremely small fraction. As ξ passes on from 1 to 2, h declines progressively [in accordance with his Eq. (3) which is valid at all nonintegral values of ξ], and at $\xi=2$ it suddenly becomes even smaller; and so on. Hence the values of h given in his Table I ($q=10$), being calculated from his erroneous equation (4), are of much too large an order at all ξ 's greater than unity; and thus the unrealistic predictions of his formulas are not surprising.

¹ P. M. Davidson, Brit. J. Appl. Phys. 4, 170 (1953).

² P. M. Davidson, Phys. Rev. 99, 1072 (1955).

³ P. L. Auer, Phys. Rev. 98, 320 (1955).

⁴ P. L. Auer, Phys. Rev. 101, 1243 (1956).

Hyperfine Structure of the Metastable State of Singly Ionized He³⁺

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THE nuclear and electrodynamic information that can be obtained from the study of atomic hyperfine structure is limited in part by the accuracy of atomic wave functions. In the case of He³ and possibly for other low- Z nuclei, this limitation can be avoided by ionizing the atom and thereby forming a hydrogen-like system whose wave function is known exactly. We describe an ion beam method for determining the hyperfine structure of the $2S_{\frac{1}{2}}$ state of ionized He³ and report a preliminary value for this quantity.

In the present experiment He³ atoms are ionized and excited by electron bombardment. At a bombarding energy of 250 ev approximately one percent of the ions is in the metastable $2S_{\frac{1}{2}}$ state. The ions are drawn out of the ion source and accelerated to 20 ev. The

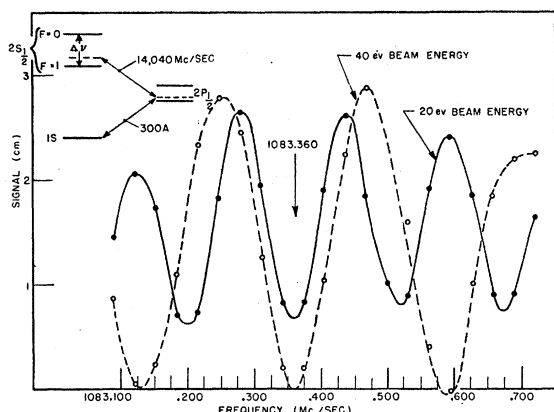


FIG. 1. The energy levels of ionized He^3 and the photoelectric current as a function of frequency for two values of the ion beam energy.

energy spread of the beam so formed is less than 1.0 eV. After passing through a lens, the ion beam traverses a first state selector, a wave guide driven at 13 350 Mc/sec. This frequency is slightly less than the separation between the lower ($F=1$) hyperfine level of the $2S_{3/2}$ state, and the upper ($F=0$) hyperfine level of the $2P_{3/2}$ state. (See insert in Fig. 1.) The rf field preferentially induces transitions from the $2S_{3/2}(F=1)$ level to the $2P_{3/2}$ level; the ions that have made the transition decay almost immediately to the ground state. As the beam leaves the wave guide there are thus more ions in the $F=0$ level than in the $F=1$ level of the $2S_{3/2}$ state.

The beam then traverses a Ramsey¹ type rf transition region over whose entire length a magnetic field of about 0.85 gauss exists. Magnetic dipole transitions are induced between the $2S_{3/2}(F=0)$ level and the $2S_{3/2}(F=1)$ level. At resonance and at optimum rf power essentially all of the excess ions in the $F=0$ state make the transition to the $F=1$ state, since the ion beam is nearly monoenergetic. The beam then passes through a second lens and traverses the second state selector, another wave guide driven at 13 350 Mc/sec. Again ions are preferentially induced to make transitions from the $2S_{3/2}(F=1)$ state to the $2P_{3/2}$ state, and they quickly decay to the ground state with the emission of 303.8 Å photons. The photons are detected photoelectrically. The magnitude of the current depends on the number of hyperfine transitions that have occurred in the Ramsey transition region. Since the beam is nearly monoenergetic, the Ramsey resonance curve exhibits a large number of fringes of almost equal intensity. The central fringe is identified by observing the pattern at two different beam velocities (see Fig. 1). The two individual transition regions of the Ramsey system are out of phase by 180° ; thus a minimum of the pattern corresponds to resonance. If the phase difference is not exactly 180° there will be a shift in the apparent value of the resonant frequency as the beam velocity is varied. By making a known

change in beam velocity this effect may be used to obtain a phase-shift correction. From five observed resonance curves at beam energies of 10, 20, and 40 eV the value of the hyperfine splitting is

$$\Delta\nu_{\text{obs}} = 1083.360 \pm 0.020 \text{ Mc/sec.} \quad (1)$$

This result includes a 2.5-kc/sec Zeeman correction and a small phase-shift correction. The uncertainty is many times the statistical error and is intended to allow for presently undetermined systematic effects such as static and rf Stark shifts. A theoretical estimate of $\Delta\nu_P(\text{He}^{3+}, 2S)$ for a point nucleus is given by the following formula

$$\begin{aligned} \Delta\nu_P(\text{He}^{3+}, 2S) = & \left[\frac{16}{3} \alpha^2 R_\infty c \left(\frac{\mu_{\text{He}^3}}{\mu_{\text{H}^1}} \right) \left(\frac{\mu_{\text{H}^1}}{\mu_S} \right) \left(\frac{\mu_S}{\mu_0} \right) \right] \\ & \times \left[1 - 3 \frac{m}{M_{\text{He}^3}} \right] \left[1 + \frac{17}{2} \alpha^2 \right] \\ & \times \left[1 + \frac{\alpha}{2\pi} - 2.973 \frac{\alpha^2}{\pi^2} \right] \\ & \times \left[1 - 2\alpha^2 \left(\frac{5}{2} - \ln 2 \right) \right]. \quad (2) \end{aligned}$$

Here the first quantity in brackets is obtained from the simple Fermi formula. The ratio of the He^3 nuclear moment to the Bohr magneton is most precisely determined by the product of 3 factors (in parentheses): the ratio of the nuclear moments of He^3 and H^1 ,² the ratio of the H^1 nuclear moment to the spin moment of the electron,³ and the theoretical value for the ratio of the electron spin moment to the Bohr magneton.⁴ The second and third quantities in brackets are respectively the Breit-Meyerott reduced-mass correction⁵ and the Breit relativistic correction.⁶ The fourth is the anomalous-moment correction for the electron and the fifth is the second-order radiative correction.⁷ Radiative corrections of relative order $Z^2\alpha^3$, $Z\alpha^3$, and α^3 are believed to exist; however these have not yet been evaluated. This formula allows for the finite nuclear mass only through the reduced-mass correction, but it does not include radiative-mass corrections⁸⁻¹⁰. Using DuMond and Cohen's¹¹ values for α , c , and R_∞ and Wapstra's¹² value for the mass of He^3 , we obtain¹³

$$\Delta\nu_P = 1083.557 \pm 0.010 \text{ Mc/sec.} \quad (3)$$

The uncertainty arises mostly from that in α ; it does not include any allowance for the unevaluated theoretical terms discussed above. We introduce a quantity δ defined by

$$\Delta\nu_{\text{obs}} = \Delta\nu_P(1 - \delta), \quad (4)$$

which is similar to the usual hyperfine structure anomaly.

Using the present value for $\Delta\nu_{\text{obs}}$ and the above

value for $\Delta\nu_P$, we obtain

$$\delta = (182 \pm 22) \text{ ppm.} \quad (5)$$

The anomaly δ presumably arises from the effects of nuclear structure, nuclear interaction currents, and the unevaluated higher order radiative corrections. Sessler and Foley¹⁴ have estimated the structure and interaction current effects. Their results depend on the choice of the nuclear wave function and on the form assumed for the interaction currents. Using the Pease-Feshbach nuclear wave function, they find a structure contribution to δ of 138 ppm. They have also considered two forms for the interaction current. The first contributes 2.0 ppm to the anomaly (δ) and the second 230 ppm. As in the case of tritium,¹⁴ the present results definitely exclude the second type of interaction current.

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¹ N. F. Ramsey, Phys. Rev. **78**, 695 (1950).

² H. L. Anderson, Phys. Rev. **76**, 1460 (1949); N. F. Ramsey, Phys. Rev. **78**, 699 (1950).

³ Koenig, Prodell, and Kusch, Phys. Rev. **88**, 191 (1952).

⁴ R. Karplus and N. M. Kroll, Phys. Rev. **77**, 536 (1950).

⁵ G. Breit and R. E. Meyerott, Phys. Rev. **72**, 1023 (1947).

⁶ G. Breit, Phys. Rev. **35**, 1447 (1930).

⁷ N. M. Kroll and F. Pollack, Phys. Rev. **86**, 876 (1952).

⁸ R. Arnowitt, Phys. Rev. **92**, 1002 (1953).

⁹ W. Newcomb and E. Salpeter, Phys. Rev. **97**, 1146 (1955).

¹⁰ Moellering, Zemach, Klein, and Low, Phys. Rev. **100**, 441 (L) (1955).

¹¹ Cohen, DuMond, Layton, and Rollett, Revs. Modern Phys. **27**, 363 (1955).

¹² A. H. Wapstra, Physica **21**, 367 (1955).

¹³ This quantity can also be obtained from the observed hfs of the ground state of hydrogen. In this case there is an uncertainty in the theory of the order of 0.020 Mc/sec (see reference 10). Ignoring the uncertain term, the value so obtained is

$$\Delta\nu_P(\text{He}^{2+}, 2S) = 1083.559 \text{ Mc/sec.}$$

¹⁴ H. M. Foley and A. M. Sessler, Phys. Rev. **98**, 6 (1955).

Effect of Temperature on the Spectral Distribution of Blue Emission Bands of ZnS:I and ZnS:Cu:I Phosphors

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IT is known that blue photoluminescence can be obtained from ZnS phosphors prepared with halide (X), as well as from ZnS phosphors prepared with high (0.1%) proportions of Cu, together with X. Some investigators¹⁻³ have felt that the blue emission from ZnS:Cu:X is due to the formation of new types of Cu centers, while others^{4,5} have believed this blue emission to be identical in origin with the blue band found in ZnS:X, which is tentatively associated with vacancies or vacancy-halide complexes.

As part of a research program on the nature of activator centers in ZnS-type phosphors, the blue emission bands were investigated with the thought that emissions arising from centers consisting of vacancies or of associated vacancies and halide would have different temperature characteristics from emissions associated with Cu. As a result, our investigations show that the blue emission bands from phosphors prepared with and without Cu can be differentiated by the effect of temperature on the spectral distribution of the emission bands.

Figure 1 shows the spectral distribution of the emission from ZnS:I and from ZnS:Cu:I, at 300°K and at 77°K, using 3650 Å excitation. The blue emission from phosphors prepared with Cu has a large *negative* temperature coefficient of the peak emission energy,

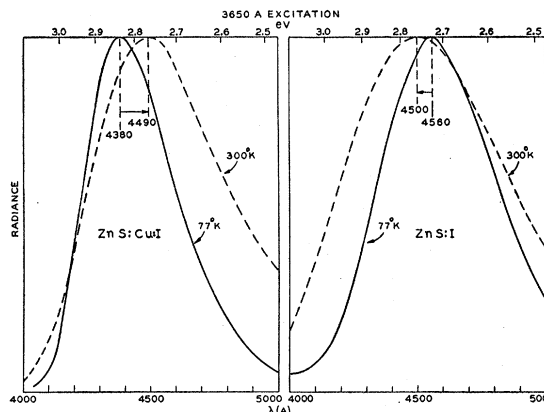


Fig. 1. Spectral distribution curves of the emission from ZnS:I and ZnS:Cu:I phosphors, at 300°K and at 77°K, using 3650 Å excitation.

while the blue emission band from phosphors prepared without Cu has a small *positive* temperature coefficient of the peak emission energy. Similar results are also found when 3125 Å photons are used for excitation.

The blue emission band from ZnS:Cu:I, therefore, should not be considered as being identical with the centers which give rise to the blue emission from ZnS:I. The utility of the temperature dependence of the spectral distribution of emission bands is further demonstrated by the finding that ZnS:Ag:X phosphors, like their Cu counterparts, also have negative temperature coefficients of the peak emission energy. In this fashion, emissions arising from vacancies (or from associated vacancy-halide complexes) can be differentiated from emissions arising from the presence of metallic activators.

¹ S. Rothschild, Trans. Faraday Soc. **42**, 635 (1946).

² N. Riehl and G. Ortman, Doklady Akad. Nauk S.S.S.R. **66**, 613 (1949).

³ F. A. Kroger and J. E. Hellingman, J. Electrochem. Soc. **93**, 156 (1948).

⁴ A. A. Cherepov and T. S. Drobolyubski, Doklady Akad. Nauk S.S.S.R. **66**, 621 (1949).

⁵ M. N. Alentsev and A. A. Cherepov, Zhur. Eksptl. i Teort. Fiz. **26**, 473 (1954).