tends to reduce the discrepancy between experiment and theory, while S and S' are kept constant. It is very apparent that $N_0 = 1{\cdot}55$ is incorrect. We find that for $N_0 = 5{\cdot}5$ and $S' = 0{\cdot}2$ the fit is reasonable. The combinations of parameters used in attempting to fit the experimental results are given in table 1.

Table 1. Parameters used to fit the high pressure data.

Density of states ratio at atmospheric pressure	Sub-band energy gap at atmospheric pressure and 295 K	Scattering parameters used in the simple theory of section 4	
$N_0 \frac{\langle 100\rangle}{\langle 111\rangle}$	ΔE_0 eV	S	S'
1·55 (Dresselhaus and Dresselhaus 1967)	0·177 ± 0·010	4·0	0·03
2·7 (Cardona and Pollak 1965)	0·180 ± 0·010	4·0	0·10
4·2	0·185 ± 0·010	4·0	0·13
5·5	0·186 ± 0·010	4·0	0·19

6. Discussion and conclusions

It becomes obvious from the curve fitting that we require N_0 to be considerably larger than has been previously theoretically advanced. The conclusion that the $\langle 100\rangle$ effective mass was considerably larger than for Si was obtained independently, but the particular value $N_0 = 5{\cdot}5$ was taken for the curve fitting since Fawcett and Paige (1971) had carried out calculations for a density of states ratio of this magnitude, and comparisons could be made.

Provided the coupling constants are known, it is possible from a determination of S to fix S' (equations 12). This enables us to compare our scattering parameters with those expected from the coupling constants used by Fawcett and Paige (1971).

For $N_0 = 5{\cdot}5$, we find $S = 4 \pm 1$ and $E_0 = 0{\cdot}186 \pm 0{\cdot}010$ eV. The coupling constants of Fawcett and Paige would give $S = 4$, which is excellent agreement. For S' however, agreement is not so good—we obtained $S' = 0{\cdot}2$, while the Fawcett and Paige theory gives $S' = 0{\cdot}5$. This discrepancy may be accounted for in the following way. The formulation developed and used to describe the L_1–Δ_1 intervalley scattering has omitted to include the phonon energy involved in the scattering event. This was noted by Nathan *et al.* (1961) but no calculations were performed. The effect of the inelastic nature of this process results in a modified expression for $\tau_g(E)$, the total relaxation time for all scattering processes in valley g

$$\frac{1}{\tau_g(E)} = A_g C'_g E^{1/2} + B_g C'_{ss} \{(E - \Delta E + \hbar\omega)^{1/2} + \exp(\hbar\omega/kT)(E - \Delta E - \hbar\omega)^{1/2}\} \times \{1 + \exp(\hbar\omega/kT)\}^{-1} \tag{13}$$

In the limit of elastic scattering the above expression reduces to equation (1) Since $\mu_H \propto \langle\tau^2(\Delta\epsilon)\rangle/\langle\tau(\Delta\epsilon)\rangle$, while $\rho \propto (\langle\tau(\Delta\epsilon)\rangle)^{-1}$ we might expect the above correction to alter the mobility fits more than the resistivity. Hence our method of fixing the point at 30 kbar in fitting the resistivity, should not incorporate a large error, and the value of $S = 4 \pm 1$ given above will reasonably allow for this. The mobility curve is particularly sensitive to S' near band cross-over, however, and here we have the largest discrepancies.

The lengthy calculations involving the use of equation (13) and the equivalent expression for $\tau_s(E)$ have not been carried out here, however, certain observations concerning the

effect of using the modified relaxation times can be made. Firstly we see that the absorption term $(E - \Delta E + \hbar\omega)^{1/2}$ makes a larger contribution to the integral derived from equation (4) than $\exp(\hbar\omega/kT)\ (E - \Delta E - \hbar\omega)^{1/2}$ where section 4 required a term $(E - \Delta E)^{1/2}$ only. Near band crossover, however, both terms become important and it is in this region that our theoretical fits disagree most with the experimental results (ie near 25 kbar). Secondly, the use of equation (13) will tend to increase the band separation at which nonequivalent intervalley scattering becomes important. The consequent reduction of the theoretical mobility before cross-over would improve the fits. Finally, the factor

$$\{1 + \exp(\hbar\omega/kT)\}^{-1}$$

may be expected to cause the calculated value of S' to be less than that defined by equation (12). The two independent results for S' are therefore converging, and perhaps a reasonable estimate is $0{\cdot}30 \pm 0{\cdot}15$.

The calculations also ignore any effect due to ionized impurity scattering. The mobility in our sample is lower than might be expected in pure Ge, and impurity scattering will probably be present. Measurements have been carried out on samples of different resistivities, however, with little variation in results. This is supported by Nathan *et al.* (1961) who found little variation at pressure for resistivities from 0·15–5 Ω cm. The results for

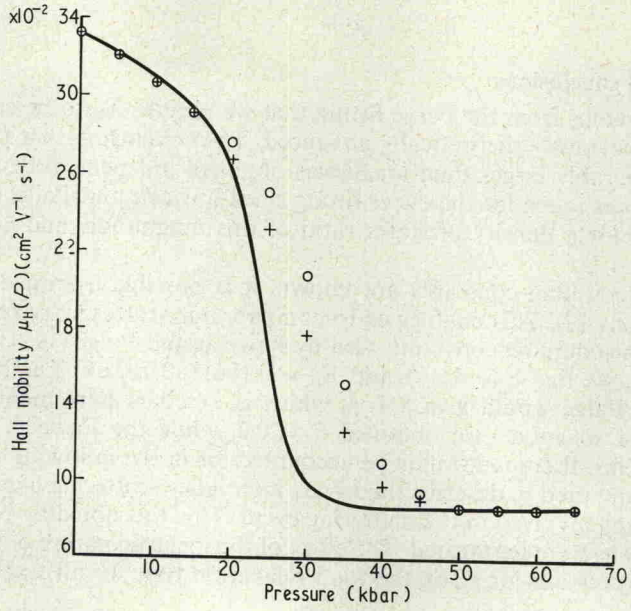

Figure 7. Theoretical fits of high pressure mobility data in n type Ge with no intervalley scattering taken into account. The largest discrepancy occurs near 30 kbar, where the true mobility is reduced by almost one half. Full curve, experimental; ○ $N = 1{\cdot}55$, $\Delta E_0 = 0{\cdot}177$ eV; + $N = 4{\cdot}2$, $\Delta E_0 = 0{\cdot}185$ eV.

our material and their samples are in excellent agreement to 30 kbar. It is reasonable to assume therefore that the mobility ratio $\mu_{H\Delta}/\mu_{H_L} = 4{\cdot}0 \pm 0{\cdot}3$, will hold for pure Ge. This implies a ⟨100⟩ mobility of $1100 \pm 150\ \text{cm}^2\ \text{V}^{-1}\ \text{s}^{-1}$ at 65 kbar. At atmospheric pressure, by analogy with the Si mobility increase with pressure and taking account of the small Ge ⟨100⟩ pressure coefficient we have $\mu_H(\Delta_1) = 1020 \pm 170\ \text{cm}^2\ \text{V}^{-1}\ \text{s}^{-1}$.

Finally, we list our conclusions. The effective mass in the Ge conduction band Δ_1 valleys is considerably greater than for Si; an estimate of 50% greater gives reasonable fits to the