The purpose of this paper is to summarize some of the research activities recently performed at the Laboratorium voor Hoge Magneto-velden at the University of Nijmegen. The scope here and unifying aspect is magnetically induced phase transitions. Here we summarize transitions in the settling velocity of paramagnetic aggregates, suppression of spin fluctuations in UA12, the phase diagram of a ferroelectric chiral smectic liquid crystal near the Lifshitz point and the transition from 3D to 2D conduction in a GaAs FET. In no way does this represent a complete review of transitions, but rather a summary of phase transitions observed at the magnet laboratory during the past year.

Magnetic Aggregation of Particles in High Fields

The stability of a colloidal suspension of particles is generally described by the Derjaguin-Landau-Verwey-Overbeek (DLVO) theory [1] where the two particle potential energy $\Phi_T$ can be expressed by the sum of potential energy $\Phi_R$ due to the electrical double layers at the solid-liquid interface, and a London-Van der Waals interaction term, $\Phi_L$. These interactions for particles of equal radii are given respectively by:

$$ \Phi_R = 4\pi \varepsilon_0 \varepsilon \varphi_0^2 \frac{\exp\left(-\frac{\gamma a}{2}\right)}{r_a} $$

(1)

and

$$ \Phi_L = -\frac{A}{6}\left[\frac{2}{r_a^2} + \frac{2}{r_a^2} + \ln\left(\frac{r_a^2 - 4}{r_a^2}\right)\right] $$

(2)

where $\varepsilon_R$ is the relative permittivity of the material of the particle, $\varepsilon_0$ the permittivity of free space, $\varphi_0$ the double layer potential, $r_a = r/a$ the normalized interparticle distance (a particle size) and $\gamma/a$ is the Debye-Hückel reciprocal double layer thickness. A potential energy barrier exists in a stable suspension which prevents the particles from approaching each other. The DLVO-theory can be extended [2] when an external magnetic field $B_0$ is turned on and the dipole-dipole interaction energy $\Phi_M$ is included. For a system of two particles, $\Phi_M$ is

$$ \Phi_M = -\frac{32\pi^2 a^3 \xi B_0^2}{9\mu_0 r_a^3}, $$

(3)
where \( \mu_0 \) is the permeability of free space and \( X \) is the net volume magnetic susceptibility of the particle. The total potential energy of interaction is given by the sum of Eqs. 1 to 3 and since \( \psi_M \) is negative, an applied magnetic field can reduce the potential energy to zero thereby allowing the particles to approach each other and form binary clusters.

In order to understand the dynamics of flocculation, a force balance equation can be constructed from the spatial derivatives of \( \phi_R \), \( \phi_L \) and \( \phi_M \). Fixing a polar coordinate system at the center of the first particle and the direction of the applied magnetic field parallel to the x-axis, the dynamics of the radial and azimuthal components of the second are respectively given by:

\[
\frac{dr_a}{dt} = -V_{Ma} \left( \frac{2 \cos^2 \alpha - 1}{r_a^3} \right) + V_{La} \left( \frac{2}{r_a^2} - \frac{1}{(r_a^2 - 4)^2} \right) - V_{Ma} \frac{r_a^2}{r_a^4} \exp \left( \frac{-r_a^2 - 2}{r_a^4} \right)
\]

(4)

\[
\frac{r_a \, d\alpha}{dt} = -V_{Ma} \frac{\sin 2\alpha}{r_a^2}
\]

(5)

Here \( V_{Ma} \), \( V_{La} \) and \( V_{Ma} \) are normalized characteristic velocity coefficients analogous to the magnetic velocity used in HGMS.

Using a fourth order Runge-Kutta technique, particle trajectories have been calculated from Eqs. 5 and 6 and are shown in Fig. 1. A paramagnetic particle initially positioned at a point \((r, \alpha)\) in the first quadrant, precesses clockwise about the “fixed” particle at the origin. If the particle is captured, it moves along the x-axis towards the point \((2a, 0)\) forming a binary pair. For particles located at a distance greater than \( r_{Ca} \), the normalized capture radius, the particle is pushed away along the x-axis.

The capture radius \( r_{Ca} \) can be estimated to a very good approximation, if \( a \) is set equal to zero in Eqs. 4 and 5. Neglecting the London-van der Waals interaction and assuming \( \tau \) is small, then the capture radius \( r_{Ca} = (2 \mu_0 / V_R) \). From the capture radius one can estimate the particle number density at which flocculation can be expected to be observed. This is of the order of

\[
\eta_0 = \left( \frac{3}{8 \cdot 4 \pi} \right) (V_R / V_M)^{3/2}
\]
Flocculation experiments were carried out on a number of commercially available mineral suspensions. The magnetic fields were produced by a water cooled Bitter magnet. Suspensions of paramagnetic materials were held in calibrated test tubes in the magnet. Field induced particle aggregation was detected by measuring the settling velocity $V_s$ of the suspended particle at different values of the applied magnetic field. This was achieved by optically measuring the rate of fall of the meniscus between sedimenting particles and clear liquid using a cathetometer or videocamera.

The settling process can be described by a simple force balance equation between the drag force and the combined influence of the gravitational and magnetic dipole forces:

$$\frac{4}{3} \pi a^3 \rho_{eff} + \frac{4}{3} \pi a^3 \chi_{eff} \frac{B_0^2}{\nu_0} = 6\pi \eta V_s$$

where $\rho_{eff}$ is the effective density of the particle in the fluid, $\chi_{eff}$ is the net volume susceptibility, $a$ is the fractional magnitude of the field gradient and $\eta$ is the viscosity of the carrier fluid. When the magnetic field exceeds $B_f$ two-particle clusters are formed which will align along the field direction, as shown in Fig. 1. With a vertical magnetic field this means that the Stoke's drag term on the RHS of Eq. 6 will remain relatively unaffected since the effective airodynamic diameter of the cluster will not be very different from that of the single particle. On the other hand, the volume dependent term on the LHS of Eq. 6 will double in size. We expect that the settling velocity plotted against $B_f$ will change discontinuously at $B_0 = B_f$ whereas its slope should approximately double.

Fig. 2.

a) Settling velocity of a 30%-vol suspension of MnO in a vertical magnetic field. The discontinuity of $V_s$ vs. $B_f$ indicates pair formation of particles. The volume susceptibility (MKSA units) is $\chi = 4.69 \times 10^{-3}$.

b) $V_s$ in a 25%-vol suspension of Mn$_2$O$_3$. $\chi$(MKSA) $= 5.39 \times 10^{-3}$.

c) $V_s$ in a 35%-vol suspension of Mn$_2$P$_2$O$_7$. $\chi$(MKSA) $= 4.72 \times 10^{-3}$. 
Measurements of settling velocity as a function of field were carried out on MnO, Mn2O3 and Mn2P2O7. Fig. 2 shows field induced binary flocculation for Mn2O3 and Mn2P2O7. In the case of MnO the slope changes by more than a factor of two; this is not clearly understood. The minimum at low fields in Mn2O3 is thought to be caused by a combination of two effects: an anisotropic magnetic susceptibility and a nonspherical particle shape. Particles rotate that their strongest magnetic axis is directed parallel to the field, thus altering their effective diameter.

Suppression of Spin Fluctuations in UA12 in High Magnetic Fields

UA12 has been known for some time to be a spin-fluctuation (SF) system. Its temperature dependence of specific heat, resistivity and magnetic susceptibility[4] fit neatly into various theories explaining these phenomena [5]. The characteristic SF temperature (Tsf) of UA12 is believed to be about 25 K. We have studied the transport properties of UA12 at low temperatures in high magnetic fields [6] where it has been argued that the application of high magnetic fields will reduce or freeze out paramagnon effects. The Zeeman splitting of the spin-up and spin-down electron states will prevent certain excitations with the reversal of spin direction to occur. The characteristic paramagnon excitations are of the order of kgTsf, with g = 2. The relationship between Hsf and Tsf (25 K) can be written as Hsf = kTsf. Hsf is then calculated to be near 19 T.

Partial freezing out of SF in high magnetic fields has been previously reported for various compounds, for example in Pd [7] Hsf is believed to be 190 T. In contrast we have achieved complete suppression of SF in UA12.

The susceptibility of UA12 [6], derived from the magnetization at 1.4 and 4.2 K, decreases with increasing field reaching a field-independent value above 20 T. At 20 and 77 K the susceptibility is found to be independent of magnetic field. Values for the differential susceptibility, see Table 1, for various field intervals as a function of temperature indicate that the high-field values at 4.2 and 1.4 K correspond to the low-field values of Brodsky and Trainor at 25 K.

Table 1 Molar susceptibilities of UA12. The units of }{M \partial \mathbf{H} \partial T} are 10^{-9} m^3/mole.

<table>
<thead>
<tr>
<th>T(K)</th>
<th>}{M \partial \mathbf{H} \partial T} (1.5 T)</th>
<th>}{M \partial \mathbf{H} \partial T} (2 - 5 T)</th>
<th>}{M \partial \mathbf{H} \partial T} (20 - 35 T)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1.4</td>
<td>-</td>
<td>56.5</td>
<td>47.3</td>
</tr>
<tr>
<td>4.2</td>
<td>55.0</td>
<td>55.4</td>
<td>47.3</td>
</tr>
<tr>
<td>20</td>
<td>-</td>
<td>47.7</td>
<td>47.7</td>
</tr>
<tr>
<td>25</td>
<td>47.3</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>50</td>
<td>46.6</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>77</td>
<td>45.5</td>
<td>45.9</td>
<td>45.9</td>
</tr>
</tbody>
</table>

a) Ref. 4 - Brodsky and Trainor.

The temperature and field dependence of the magnetic susceptibility suggest that an equivalent contribution can be suppressed either by applying a magnetic field of about 20 T or increasing the temperature to Tsf.

The temperature dependent susceptibility has also been measured from 4.2 to 270 K at 1.2 and 8.0 T. The low temperature data corresponds to an expression X = X0 (1 - T^2/Tsf). The change of the specific heat in a magnetic field can be predicted from the susceptibility through the Maxwell relation,

\[ \left( \frac{\partial S}{\partial V_{\mathbf{H}}} \right)_{T,p} = \left( \frac{\partial M}{\partial T} \right)_{H,p}. \]
On taking the temperature derivative and using $M = xH$, we find that
\[
\left( \frac{\partial \gamma}{\partial \mu^2} \right)_T = 2 \left( \frac{\partial \chi}{\partial T^2} \right)_H \rho
\]  
(8)

Since the electronic contribution to the specific heat is linear in temperature and using the $T^2$ fit to the susceptibility, we find that Eq. 8 can be written as
\[
\left( \frac{\partial \gamma}{\partial \mu^2} \right)_T = -2 \left( \frac{x_0^2}{\gamma_H^2} \right) \rho
\]  
(9)

From Table 1, $x_0$ decreases with increasing field thus the specific-heat coefficient $\gamma$ diminishes with a power law less than $H^2$. The specific heat experiments by Trainor et al. [4] are in very good agreement with the suppression of $\gamma_H$ calculated here to be $-1.25$ at 4 T.

The magnetoresistance shows a bending over at fields of about 15 T, as shown in Fig. 3. The power law is $\Delta \rho = H^n$, for fields between 0 and 15 T, $n = 1.45$, while above 15 T, $n = 1.3$. Hertel, Appel and Fay [8] calculated the magnetic field dependence of UA$_2$ where a 14% suppression of the SF contribution is predicted at 10 T. If we assume that the magnetoresistance due to paramagnons and other effects (band, $e - e$, $e - p$, etc.) are in series then the total resistivity may be written as $\Delta \rho = \Delta \rho_{SF} + \Delta \rho_{other}$. If we additionally assume that $\rho_{SF}$ is maximum at zero field (contributing 30% at $H = 0$, from an extrapolation above 15 T) then the paramagnon contribution is suppressed by 75% at 10 T.

We find that the paramagnon effects at low temperature in UA$_2$ are large and are frozen in fields above $H_{SF}$. The characteristic field and temperature are related by $g_{SF} H = g_{SF} T$, with $g = 2$. As a result of suppression of $SF$ (above $H_{SF}$) it is quite plausible that the de Haas-van Alphen effect can only be seen when all paramagnon contributions to the effective mass are frozen out. In preliminary experiments, dHvA oscillations have been observed in UA$_2$ only above 20 T [9]. In addition, these experiments suggest that field and temperature effects on the paramagnon contributions to the transport properties are closely related.

**Fig. 3** Magnetoresistivity of UA$_2$ at 4.2 K.

**Magnetic Phase Diagram of a Ferroelectric Chiral Smectic Liquid Crystal near the Lifshitz Point**

The Lifshitz point [10] is that point on a line of second-order phase transitions where an instability occurs in wave vector space. It represents a special case of a triple point between a disordered, homogeneously ordered ($R = 0$), and modulated phase ($R \neq 0$). It has been predicted [11] that a Lifshitz point exists in the $T - H$ phase diagram of chiral smectic liquid crystals if the magnetic field is applied parallel to the smectic layers.
Here the phase diagram of the chiral ferroelectric smectic liquid crystal p-decyl-oxybenzilidene-p'-amino-2-methylbutylcinnamate (DOBAMBC) is presented [12] which suggests that a Lifshitz point exists between the disordered smectic-A, the helicoidally ordered smectic-C*, and the homogeneously ordered smectic-C phases. The observed $T_c - H$ phase diagram in magnetic field up to 14.5 T is shown in Fig. 4. The A - C* line was found to be of second order whereas the C - C* transition line is of first order close to the $\lambda$ line. The discontinuity in $K$ becomes smaller as the Lifshitz point is approached along the C* - C boundary. Above $H = 8.5$ T the C - C* line changes its direction so that a reentrant smectic-C* phase was found.

The temperature and magnetic field dependence of the in-plane component of the dielectric constant $\varepsilon$ and the pitch $2n/k$ of the smectic-C* helix were measured in order to determine the phase boundaries. Simultaneous dielectric and optical measurements were performed using optically flat glass plates with transparent SnO$_2$ electrodes. The molecules were oriented parallel to the cell walls by slowly cooling (1°C/h) the system through the isotropic-smectic A transition in a magnetic field of 10 T. The C - C* phase boundary and the critical magnetic field $H_L$ for unwinding the smectic-C* helix was determined by measuring $\varepsilon = \varepsilon(H)$ at $T = \text{const}$. The sweep rate was 1.45 T/min. The A - C* boundary was determined by measuring $\varepsilon = \varepsilon(T)$ at $H = \text{const}$.

The inverse pitch $(k/n)$ must tend to zero as the Lifshitz point is approached. If not, the $k$ of the C - C* transition is discontinuous at the A - C* - C triple point. The magnetic field dependence of the wave vector of the C* helix along the $\lambda$ line has been given [2] as

$$k_2^2(H) = k_0^2(1 - H/H_L^0), \quad H < H_L$$

while $k_0 = 0$ for $H > H_L$. Fig. 5 shows $k$ measured as a function of $H$ for constant $T$. The discontinuity at the C* - C transition decreasing with increasing $T$ as the Lifshitz point is approached. The isotherm ($T_c - T = 0.2$) can be reasonably fit by Eq. 10 with $H_L = 10.2$ T. However, for $T_c - T = 0.1$, $H_L$ is larger than 14 T. The system behaves in the way predicted [2] for $T_c - T > 0.2$, however, for $T_c - T < 0.2$ the critical field for unwinding the helix strongly increases with increasing temperature.
Four experimental measurements were made on each device [14]: capacitance, resistance, Hall effect and Shubnikov-de Haas effect. The doping concentration $N_D$ could then be independently confirmed from the capacitance measurements on one hand and the SdH and Hall effect on the other. SdH oscillations for various $V_s$'s with constant substrate voltages were measured and from the periodicity of their oscillations, the Fermi energy is readily determined from

$$n_L = \frac{m^*E_F}{\hbar eB} - \frac{1}{2}$$

(11)

where $n_L$ is the Landau level index. For the device with a doping concentration of $N_D = 2 \times 10^{17} \text{ cm}^{-3}$, the Fermi energy is plotted in Fig. 6. From the oscillatory part of the magnetoresistance

$$\cos \left( \frac{2\pi(E_F - E_n)}{\hbar} - \alpha \right)$$

(12)
each sub-band contributes a periodicity dependent on $\Delta E_n = E_F - E_n$. The theoretical oscillation periodicity obtained from Eq. 12 was used to calculate the sub-band energies in Fig. 6 [14]. As $t$ is decreased towards zero it will eventually become less than the mean donor separation $\bar{r}$ which is 174 Å. A transition from metallic behaviour was observed via the quenching of SdH oscillations for values of $t < \bar{r}$. It can be seen, in Fig. 6, that there is a reasonable agreement between the experimental points and the theoretical predictions for $E_F$. The fall in $E_F$ as the conducting electron gas in a GaAs FET is confined is confirmed resulting from a 3D to 2D transition.

Fig. 6 Level quantization and $E_F$ as a function of channel thickness ($t$) in a $N_D = 2 \times 10^{17}$ cm$^{-3}$ device. Solid curves are WKB level separation (see ref. 14) $\Delta E_n$ and the chain curve to the corresponding $E_F$. The straight horizontal broken line is $E_F$ (3D) for a uniform non-interacting gas. Measured $E_F$ values are presented as full circles with error bars.

Here we have presented a variety of magnetically induced phase transitions measured at the High Field Magnet Laboratory of the University during this past year. The extent to which magnetic fields help us understand the properties of matter remains significant.

Acknowledgments

We would like to thank Klaas van Hulst, Henk Muileman and Jos Rook of the Nijmegen Magnet Laboratory for their valuable assistance and technical support of this research. Part of this work has been supported by the "Stichting voor Fundamenteel Onderzoek der Materie" (FOM) with financial support of the "Nederlandse Organisatie voor Zuiver Wetenschappelijk Onderzoek" (ZWO).

References